Why the Quantum Must Yield to Gravity

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ABSTRACT

After providing an extensive overview of the conceptual elements – such as Einstein's 'hole argument' – that underpin Penrose's proposal for gravitationally induced quantum state reduction, the proposal is constructively criticised. Penrose has suggested a mechanism for objective reduction of quantum states with postulated collapse time $\tau = \hbar/\Delta E$, where ΔE is an ill-definedness in the gravitational self-energy stemming from the profound conflict between the principles of superposition and general covariance. Here it is argued that, even if Penrose's overall conceptual scheme for the breakdown of quantum mechanics is unreservedly accepted, his formula for the collapse time of superpositions reduces to $\tau \to \infty \ (\Delta E \to 0)$ in the strictly Newtonian regime, which is the domain of his proposed experiment to corroborate the effect. A suggestion is made to rectify this situation. In particular, recognising the cogency of Penrose's reasoning in the domain of full 'quantum gravity', it is demonstrated that an appropriate experiment which could in principle corroborate his argued 'macroscopic' breakdown of superpositions is not the one involving non-rotating mass distributions as he has suggested, but a Leggett-type SQUID or BEC experiment involving superposed mass distributions in relative rotation. The demonstration thereby brings out one of the distinctive characteristics of Penrose's scheme, rendering it empirically distinguishable from other state reduction theories involving gravity. As an aside, a new geometrical measure of gravity-induced deviation from quantum mechanics à la Penrose is proposed, but now for the canonical commutation relation $[Q, P] = i\hbar$.

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1. Introduction – From Schrödinger's Cat to Penrose's 'OR':

Quantum mechanics – one of our most fundamental and successful theories – is infested with a range of deep philosophical difficulties collectively known as the measurement problem (Schrödinger 1935, Shimony 1963, Wheeler and Zurek 1983, Bell 1990). In a nutshell the problem may be stated as follows: If the orthodox formulation of quantum theory – which in general allows attributions of only objectively indefinite properties or potentialities (Heisenberg 1958) to physical objects – is interpreted in compliance with what is usually referred to as scientific realism, then one is faced with an irreconcilable incompatibility between the nonnegotiable linearity of quantum dynamics – which governs evolution of the network of potentialities – and the apparent definite or actual properties of the physical objects of our 'macroscopic' world. Moreover, to date no epistemic explanation of these potentialities (e.g., in terms of 'hidden variables') has been completely successful (Shimony 1989). Thus, on the one hand there is overwhelming experimental evidence in favour of the quantum mechanical potentialities, supporting the view that they comprise a novel (i.e., classically uncharted) metaphysical modality of Nature situated between logical possibility and actuality (Shimony 1978, 1993b, 140-162 and 310-322), and on the other hand there is phenomenologically compelling proliferation of actualities in our everyday world, including even in the microbiological domain. The problem then is that a universally agreeable mechanism for transition between these two ontologically very different modalities i.e., transition from the multiplicity of potentialities to various specific actualities – is completely missing. As delineated, this is clearly a very serious physical problem. What is more, as exemplified by Shimony (1993a, 56), the lack of a clear understanding of this apparent transition in the world is also quite a 'dark cloud' for any reasonable program of scientific realism.

Not surprisingly, there exists a vast number of proposed solutions to the measurement problem in the literature (Christian 1996), some of which – the Copenhagen interpretation (Bohr 1935) for example – being almost congenital to quantum mechanics. Among these proposed solutions there exists a somewhat dissident yet respectable tradition of ideas – going all the way back to Feynman's pioneering thoughts on the subject as early as in mid-fifties (Feynman 1957) – on a possible gravitational resolution of the problem. The basic tenet of these proposals can hardly be better motivated than in Feynman's own words. In his Lectures on Gravitation (Feynman 1995, 12-13) he devotes a whole section to the issue, entitled "On the philosophical problems in quantizing macroscopic objects", and contemplates on a possible breakdown of quantum mechanics:

"...I would like to suggest that it is possible that quantum mechanics fails at large distances and for large objects. Now, mind you, I do not say that I think that quantum mechanics does fail at large distances, I only say that it is not inconsistent with what we do know. If this failure of quantum mechanics is connected with gravity, we might speculatively expect this to happen for masses such that $GM^2/\hbar c=1$, of M near 10^{-5} grams, which corresponds to some 10^{18} particles."

Indeed, if quantum mechanics does fail near the Planck mass, as that is the mass scale Feynman is referring to here, then – at last – we can put the annoying problem of measurement to its final rest (see Figure 1 for the meanings of the constants G, \hbar , and c). The judiciously employed tool in practice, the infamous postulate often referred to as the reduction of quantum state – which in orthodox formulations of the theory is taken as one of the unexplained basic postulates to resolve the tension between the linearity of quantum dynamics and the plethora of physical objects with apparent definite properties – may then be understood as an objective physical phenomenon; i.e., one affording an ontological as opposed to epistemological understanding. From the physical viewpoint such a resolution of the measurement problem would be quite satisfactory, since it would render the proliferation of diverse philosophical opinions on the matter to nothing more than a curious episode in the history of physics. For those who are not lured by pseudo-solutions such as the 'decohering histories' approaches (Kent 1997) and/or 'many worlds' approaches (Kent 1990), a resolution of the issue by 'objective reduction' ('OR', to use Penrose's ingenious pun) comes across as a very attractive option, provided of course that that is indeed the path Nature has chosen to follow (see also Christian 1999a).

Motivated by Feynman's inspiring words quoted above, there have been several concrete theoretical proposals of varied sophistication and predilections on how the breakdown of quantum mechanics might come about such that quantum superpositions are maintained only for 'small enough' objects, whereas reduction of

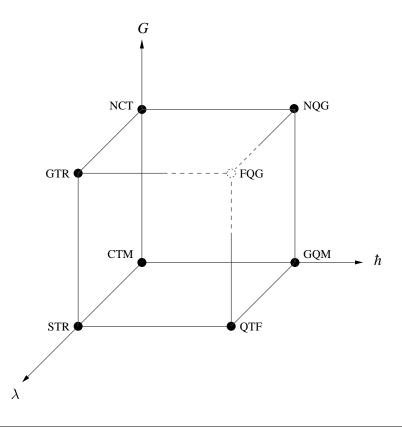


Figure 1: The great dimensional monolith of physics indicating the fundamental role played by the three universal constants G (the Newton's gravitational constant), \hbar (the Planck's constant of quanta divided by 2π), and $\lambda \equiv 1/c$ (the 'causality constant' (Ehlers 1981), where c is the absolute upper bound on the speed of causal influences) in various basic theories. These theories, appearing at the eight vertices of the cube, are: CTM = Classical Theory of Mechanics, STR = Special Theory of Relativity, GTR = General Theory of Relativity, NCT = Newton-Cartan Theory, NQG = Newton-Cartan Quantum Gravity (Christian 1997), GQM = Galilean-relativistic Quantum Mechanics, QTF = Quantum Theory of (relativistic) Fields, and FQG = the elusive Full-blown Quantum Gravity. Note that FQG must reduce to QTF, GTR, and NQG in the respective limits $G \to 0$, $\hbar \to 0$, and $\lambda \to 0$ (Kuchař 1980, Christian 1997, and Penrose 1997, 90-92).

the quantum state is objectively induced by gravity for 'sufficiently large' objects (Károlyházy 1966, Komar 1969, Kibble 1981, Károlyházy et al. 1986, Diósi 1984, 1987, 1989, Ellis et al. 1989, Ghirardi et al. 1990, Christian 1994, Percival 1995, Jones 1995, Pearle and Squires 1996, Frenkel 1997, Fivel 1997). Unfortunately, most of these proposals employ dubious or ad hoc notions such as 'quantum fluctuations of spacetime' (e.g., Percival 1995) and/or 'spontaneous localization of the wavefunction' (e.g., Ghirardi et al. 1990). Since the final 'theory of everything' or 'quantum gravity' is quite far from enjoying any concrete realization (Rovelli 1998), such crude notions cannot be relied upon when discussing issues as fundamental as the measurement problem. In fact, these notions are not just unreliable, but, without the context of a consistent 'quantum theory of gravity', they are also quite meaningless. For this reason, in this essay I shall concentrate exclusively on Penrose's proposal of quantum state reduction (1979, 1981, 1984, 1986, 1987, 1989, 367-371, 1993, 1994a, 1994b, 339-346, 1996, 1997, 1998), since his is a minimalist approach in which he refrains from unnecessarily employing any ill-understood (or oxymoronic) notions such as 'quantum fluctuations of spacetime'. Rather, he argues from the first principles, exploiting the profound and fundamental conflict (all too familiar to anyone who has attempted to 'quantize gravity' – sometimes in the guise of the so-called 'problem of time'

(Kuchař 1991, 1992, Isham 1993, Belot and Earman 1999)) between the principle of general covariance of general relativity and the principle of superposition of quantum mechanics, to deduce a heuristic mechanism of gravity-induced quantum state reduction. Stated differently, instead of prematurely proposing a crude theory of quantum state reduction, he merely provides a rationale for the mass scale at which quantum mechanics must give way to gravitational effects, and hence to a superior theory.

Let me emphasize further that the motivations based on rather contentious conceptual issues inwrought in the measurement problem are not an essential prerequisite to Penrose's proposal for the breakdown of quantum superpositions at a 'macroscopic' scale. Instead, his proposal can be viewed as a strategy not only to tackle the profound tension between the foundational principles of our two most fundamental physical theories – general relativity and quantum mechanics, but also to simultaneously provide a possible window of opportunity to go beyond the confining principles of these two great theories in order to arrive at even greater enveloping 'final' theory (Penrose 1984, Christian 1999b). Such a final theory, which presumably would neither be purely quantal nor purely gravitational but fundamentally different and superior, would then have to reduce to quantum mechanics and general relativity, respectively, in some appropriate approximations, as depicted in Figure 1. Clearly, unlike the lopsided orthodox approaches towards a putative 'quantum theory of gravity' (Rovelli 1998), this is a fairly 'evenhanded' approach – as Penrose himself often puts it. For, in the orthodox approaches, quantum superpositions are indeed presumed to be sacrosanct at all physical scales, but only at a very high price of some radical compromises with Einstein's theory of gravity (e.g., at a price of having to fix both the topological and differential structures of spacetime a priori, as in the 'loop quantum gravity' program (Rovelli 1998), or – even worse – at a price of having to assume some non-dynamical causal structure as a fixed arena for dynamical processes, as in the currently voguish 'M-theory' program (Banks 1998a, b, Polchinski 1998, Sen 1998), either of the compromises being anathematic to the very essence of general relativity (Einstein 1994, 155, Stachel 1994, Isham 1994, Sorkin 1997)).

In passing, let me also point out another significant feature of Penrose's proposal which, from a certain philosophical perspective (namely the 'process' perspective (Whitehead 1929)), puts it in a class of very attractive proposals. Unlike some other approaches to the philosophical problems of quantum theory, his approach (and for that matter almost all approaches appealing to the 'objective reduction') implicitly takes temporal transience in the world – the incessant fading away of the specious present into the indubitable past – not as a merely phenomenological appearance, but as a bona fide ontological attribute of the world, in a manner, for example, espoused by Shimony (1998). For, clearly, any gravity-induced or other intrinsic mechanism, which purports to actualize – as a real physical process – a genuine multiplicity of quantum mechanical potentialities to a specific choice among them, evidently captures transiency, and thereby not only goes beyond the symmetric temporality of quantum theory, but also acknowledges the temporal transience as a fundamental and objective attribute of the physical world (Shimony 1998) (for anticipatory views on 'becoming' along this line, see also (Eddington 1929, Bondi 1952, Reichenbach 1956, Whitrow 1961)). A possibility of an empirical test confirming the objectivity of this facet of the world via Penrose's approach is by itself sufficient for me to endorse his efforts wholeheartedly. But his approach has even more to offer. It is generally believed that the classical general relativistic notion of spacetime is meaningful only at scales well above the Planck regime, and that near the Planck scale the usual classical structure of spacetime emerges purely phenomenologically via a phase transition or symmetry breaking phenomenon (Isham 1994). Accordingly, one may incline to think that "the concept of 'spacetime' is not a fundamental one at all, but only something that applies in a 'phenomenological' sense when the universe is not probed too closely" (Isham 1997). However, if the emergence of spacetime near the Planck scale is a byproduct of the actualization of quantum mechanical potentialities – via Penrose's or any related mechanism, then the general relativistic spacetime, along with its distinctively dynamical causal structure, comes into being not as a coarse-grained phenomenological construct, but as a genuine ontological attribute of the world, in close analogy with the special case of temporal transience. In other words, such an ontological coming into being of spacetime near the Planck scale would capture the 'becoming' not merely as temporal transience, which is a rather 'Newtonian' notion, but as a much wider, dynamical, spatio-temporal sense parallelling general relativity. (This will become clearer in section 4 below where I discuss Penrose's mechanism, which is tailor-made to actualize specific spacetime geometries out of 'superpositions' of such geometries). This gratifying possibility leaves no shred of doubt that the idea of 'objective reduction' in general, and its variant proposed by Penrose in particular, is worth investigating seriously, both theoretically and experimentally¹.

Since the principle of general covariance is at the heart of Penrose's proposal, I begin in the next section with a closer look at the physical meaning of this fundamental principle, drawing lessons from Einstein's struggle to come to terms with it by finding a resolution of his famous 'hole argument' (1914). Even the reader fairly familiar with this episode in the history of general relativity is urged to go through the discussion offered here, since the subtleties of the principle of general covariance provides the basis for both Penrose's central thesis as well as my own partial criticism of it. Next, after highlighting the inadequacies of the orthodox quantum measurement theory in section 3, I review Penrose's proposal in greater detail in section 4, with a special attention to the experiment he has proposed to corroborate his quantitative prediction of the breakdown of quantum mechanics near a specific mass scale (subsection 4.5). (As an aside, I also propose a new geometrical measure of gravity-induced deviation from quantum mechanics in subsection 4.4.) Since Penrose's proposed experiment is entirely within the nonrelativistic domain, in the subsequent subsection, 5.1, I provide an orthodox analysis of it strictly within this domain, thereby setting the venue for my partial criticism of his proposal in subsection 5.2. The main conclusion here is that, since there remains no residue of the conflict between the principles of superposition and general covariance in the strictly-Newtonian limit (and this happens to be a rather subtle limit), Penrose's formula for the 'decay-time' of quantum superpositions produces triviality in this limit, retaining the standard quantum coherence intact. Finally, in subsection 5.3, before making some concluding remarks in section 6, I suggest that an appropriate experiment which could in principle corroborate Penrose's predicted effect is not the one he has proposed, but a Leggett-type SQUID (Superconducting QUantum Interference Devise) or BEC (Bose-Einstein Condensate) experiment involving superpositions of mass distributions in relative rotation. As a bonus, this latter analysis brings out one of the distinctive features of Penrose's scheme, rendering it empirically distinguishable from all of the other (ad hoc) quantum state reduction theories involving gravity (e.g., Ghirardi et al. 1990).

2. How the Spatio-temporal Events Lost Their Individuality:

Between 1913 and 1915 Einstein (1914) put forward several versions of an argument, later termed by him the 'hole argument' ('Lochbetrachtung'), to reject what is known as the principle of general covariance, which he himself had elevated earlier as a criterion for selecting the field equations of a theory of gravitation he was in a process of constructing. It is only after two years of struggle to arrive at the correct field equations with no avail that he was led to reconsider general covariance, despite the hole argument, and realized the full significance and potency of the principle it enjoys today. In particular – and this is also of utmost significance for our purposes here – he realized that the hole argument and the principle of general covariance can peacefully coexist if and only if the mathematical individuation of the points of a spacetime manifold is physically meaningless. In other words, he realized that a bare spacetime manifold without some 'individuating field' (Stachel 1993) such as a specific metric tensor field defined on it is a highly fictitious mathematical entity without any direct physical content.

Although the physical meaninglessness of a mathematical individuation of spacetime points – as a result of general covariance – is central to Penrose's proposal of quantum state reduction, he does not invoke the historical episode of hole argument to motivate this nontrivial aspect of the principle. And justifiably so. After all, the non-triviality of the principle of general covariance (i.e., the freedom under active diffeomorphisms of spacetime) is one of the first things one learns about while learning general relativity. For example, Hawking and Ellis (1973) begin their seminal treatise on the large scale structure of spacetime by simply taking a mathematical model of spacetime to be the entire equivalence class of copies of a 4-manifold, equipped, respectively, with Lorentzian metric fields related by active diffeomorphisms of the manifold, without even mentioning the hole argument. However, as we shall see, it is the hole argument – an argument

¹ It should be noted that Penrose's views on 'becoming' are rather different from the stance I have taken here (1979, 1989, 1994b). In the rest of this essay I have tried to remain as faithful to his writings as possible. For recent discussions and further references on 'becoming', other than the paper by Shimony cited above, see (Zeilicovici 1986), (Saunders 1996), and (Magnon 1997).

capable of misleading even Einstein for two years – that demands such an identification in the first place. Therefore, and especially considering the great deal of persistent confusion surrounding the physical meaning of the principle of general covariance in the literature (Norton 1993), for our purposes it would be worthwhile to take a closer look at the hole argument, and thereby appreciate what is at the heart of Penrose's proposal of quantum state reduction. For more details on the physical meaning of general covariance the reader is referred to Stachel's incisive analysis (1993) of it in the modern differential geometric language; it is the general viewpoint espoused in this reference that I shall be mostly following here (but see also (Rovelli 1991) and section 6 of (Anandan 1997) for somewhat analogous viewpoints).

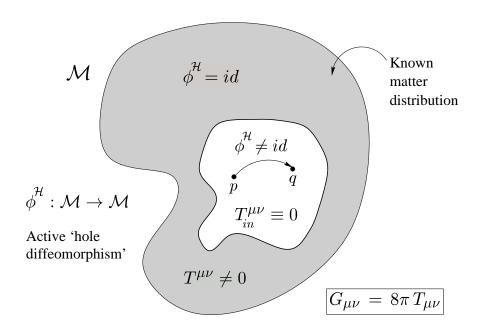


Figure 2: Einstein's hole argument: If the field equations of a gravitational theory are generally covariant, then, inside a matter-free region of some known matter distribution, they appear to generate infinite number of inequivalent solutions related by active diffeomorphisms of the underlying spacetime manifold.

Without further ado, here is Einstein's hole argument: As depicted in Figure 2, suppose that the matter distribution encoded in a stress-energy tensor $T^{\mu\nu}$ is precisely known everywhere on a spacetime \mathcal{M} outside of some hole $\mathcal{H} \subset \mathcal{M}$ – i.e., outside of an open subspace of the manifold \mathcal{M} . (Throughout this essay I shall be using Penrose's abstract index notation (Wald 1984).) Further, let there be no physical structure defined within \mathcal{H} except a gravitational field represented by a Lorentzian metric tensor field $g^{\mu\nu}$; i.e., let the stress-energy vanish identically inside the hole: $T_{in}^{\mu\nu} \equiv 0$. Now suppose that the field equations of the gravitational theory under consideration are generally covariant. By definition, this means that if a tensor field \mathcal{X} on the manifold \mathcal{M} is a solution of the set of field equations, then the pushed-forward tensor field $\phi_*\mathcal{X}$ of \mathcal{X} is also a solution of the same set of equations for any active diffeomorphism $\phi: \mathcal{M} \to \mathcal{M}$ of the manifold \mathcal{M} onto itself. The set of such diffeomorphisms of \mathcal{M} forms a group, which is usually denoted by Diff(\mathcal{M}). It is crucially important here to distinguish between this genuine group Diff(\mathcal{M}) of global diffeomorphisms of \mathcal{M} and the pseudo-group of transformations between overlapping pairs of local coordinate charts. The elements of the latter group are sometimes referred to as passive diffeomorphisms because they can only produce trivial transformations by merely relabelling or renaming the points of a manifold. Admittance of only tensorial objects on \mathcal{M} in any spacetime theory is sufficient to guarantee compatibility with this

pseudo-group of passive diffeomorphisms. On the other hand, the elements of the genuine group $Diff(\mathcal{M})$ of active diffeomorphisms are smooth homeomorphisms $\phi: \mathcal{M} \to \mathcal{M}$ that can literally take each point p of \mathcal{M} into some other point $q := \phi(p)$ of \mathcal{M} and thereby deform, for example, a doughnut shaped manifold into its coffee-mug shaped copy (Nakahara 1990, 54). Returning to the definition of general covariance, if a metric tensor field $q^{\mu\nu}(x)$ is a solution of the generally covariant field equations at any point x of \mathcal{M} in an adapted local coordinate system, then so is the corresponding pushed-forward tensor field $(\phi_* g)^{\mu\nu}(x)$ at the same point x in the same coordinate system. Note that, in general, $g^{\mu\nu}$ and $(\phi_*g)^{\mu\nu}$ will be functionally different from each other in a given coordinate system; i.e., the components of $(\phi_* g)^{\mu\nu}$ will involve different functions of the coordinates compared to those of $g^{\mu\nu}$. Now, since a choice of $\phi \in \text{Diff}(\mathcal{M})$ is by definition arbitrary, nothing prevents us from choosing a smooth $\phi^{\mathcal{H}}$ – a 'hole diffeomorphism' – which reduces to $\phi^{\mathcal{H}} = id$ (i.e., identity) everywhere outside and on the boundary of the hole \mathcal{H} , but remains $\phi^{\mathcal{H}} \neq id$ within \mathcal{H} . Such a choice, owing to the fact $T^{\mu\nu} \equiv 0$ within the hole, implies that the action of $\phi^{\mathcal{H}}$ will not affect the stress-energy tensor anywhere: $(\phi_*^{\mathcal{H}}T)^{\mu\nu} = T^{\mu\nu}$ everywhere, both inside and outside of \mathcal{H} . On the other hand, applied to the metric tensor $g^{\mu\nu}$, $\phi^{\mathcal{H}}$ will of course produce a new solution of the field equations according to the above definition of general covariance, although outside of \mathcal{H} this new solution will remain identical to the old solution. The apparent difficulty, then, is that, even though $T^{\mu\nu}$ remains unchanged, our choice of the hole diffeomorphism ϕ^{μ} allows us to change the solution $g^{\mu\nu}$ inside the hole as non-trivially as we do not like, in a blatant violation of the physically natural uniqueness requirement, which states that the distribution of stress-energy specified by the tensor $T^{\mu\nu}$ should uniquely determine the metric tensor $g^{\mu\nu}$ representing the gravitational field. Indeed, under the diffeomorphism $\phi^{\mathcal{H}}$, identical matter fields $T^{\mu\nu}$ seam to lead to non-trivially different gravitational fields inside the hole, such as $g^{\mu\nu}$ and $(\phi^{\mathcal{H}}_{*}g)^{\mu\nu}$, since $\phi^{\mathcal{H}}$ is not an identity there. What is worse, even though nothing has been allowed to change outside or on the boundary of the hole \mathcal{H} , nothing seams to prevent the gravitational field $(\phi^{\mathcal{H}}_{*}g)^{\mu\nu}$ from being completely different for each one of the *infinitely many* inequivalent diffeomorphisms $\phi^{\mathcal{H}} \in \mathrm{Diff}(\mathcal{M})$ that can be carried out inside \mathcal{H} .

As mentioned above, Einstein's initial reaction to this dilemma was to abandon general covariance for the sake of uniqueness requirement, and he maintained this position for over two years. Of course, to a modern general relativist a resolution of the apparent problem is quite obvious: The tacit assumption in the hole argument that the mathematically different tensor fields $g^{\mu\nu}$ and $(\phi_*^{\mathcal{H}}g)^{\mu\nu}$ are also physically different – i.e., correspond to different physical realities – is clearly not justified. The two expressions $g^{\mu\nu}$ and $(\phi_*^{\mathcal{H}}g)^{\mu\nu}$ may no matter how non-trivially differ mathematically, they must represent one and the same gravitational field physically. Thus, as Wald puts it (1984), "diffeomorphisms comprise [nothing but a] gauge freedom of any theory formulated in terms of tensor fields on a spacetime manifold". Accordingly, in formal analogy with the familiar gauge freedom of the gauge field theories, modern general relativists take a gravitational field to physically correspond to an entire equivalence class of metric tensor fields, related by arbitrary diffeomorphisms of the spacetime manifold, and not just to one of the members of this class.

The analogy with the gauge freedom of the gauge field theories, however, has only a limited appeal when it comes to general relativity. To see the difference, recall, for example, that electromagnetic gauge transformations – the prototype of all gauge transformations – occur at a fixed spacetime point: The vector potential $A_{\mu}(x)$ defined at a point x of \mathcal{M} is physically equivalent to the vector potential $A_{\mu}(x) + \partial_{\mu} f(x)$ defined at the same point x of \mathcal{M} , for all scalar functions f(x). Although mathematically different, both $A_{\mu}(x)$ and $A_{\mu}(x) + \partial_{\mu} f(x)$ correspond to one and the same physical electromagnetic field configuration $F_{\mu\nu}(x)$, which again depends locally on the same point x of \mathcal{M} . On the other hand, as stressed above, in general relativity diffeomorphisms $\phi \in \text{Diff}(\mathcal{M})$ map one spacetime point, say p, to another spacetime point, say $q := \phi(p)$. Therefore, if the tensor fields $g^{\mu\nu}(p)$ and $(\phi_* g)^{\mu\nu}(q)$ are to be identified as representing one and the same gravitational field configuration, implying that they cannot be physically distinguishable by any means, then the two points p and p must also be physically indistinguishable, and, consequently, they must renounce their individuality. For, if the points of \mathcal{M} did possess any ontologically significant individual identity of their own, then a point p of \mathcal{M} could be set apart from a point p of \mathcal{M} , and that would be sufficient to distinguish the quantity $g^{\mu\nu}(p)$ from the quantity $(\phi_* g)^{\mu\nu}(q)$, contradicting the initial assertion.

As Einstein eventually realized, the conclusion is inescapable: The points of a spacetime manifold \mathcal{M} have no direct ontological significance. A point in a bare spacetime manifold is not distinguishable from any other point - and, indeed, does not even become a point (i.e, an event) with physical meaning - unless and until a specific metric tensor field is dynamically determined on the manifold. In fact, in general relativity a bare manifold not only lacks this local property, but the entire qlobal topological structure of spacetime is also determined only a posteriori via a metric tensor field (Einstein 1994, Stachel 1994, Isham 1994, Sorkin 1997). Since a dynamical metric tensor field on a manifold dynamizes the underlying topology of the manifold, in general relativity the topology of spacetime is also not an absolute element that 'affects without being affected'. Thus, strictly speaking, the bare manifold does not even become 'spacetime' with physical meaning until both the global and local spatio-temporal structures are dynamically determined along with a metric. Further, since spacetime points aguire their individuality in no other way but as a byproduct of a solution of Einstein's field equations, in general relativity 'here' and 'now' cannot be part of a physical question, but can only be part of the answer to a question, as Stachel so aptly puts it (1994). The concepts 'here' and 'now' - and hence the entire notion of local causality - aquire ontological meaning only a posteriori, as a part of the answer to a physical question. Anticipating the issue discussed in the section 4 below, this state of affairs is in sharp contrast to what one can ask in quantum theory, which due to its axiomatically non-dynamical causal structure – allows 'here' and 'now' to be part of a question. Indeed, in quantum mechanics, as we shall see, a priori individuation of spatio-temporal events is an essential prerequisite to any meaningful notion of time-evolution.

At a risk of repetition, let me recapitulate the central point of this section in a single sentence:

In Einstein's theory of gravity, general covariance – i.e., invariance of physical laws under the action of the group $Diff(\mathcal{M})$ of active diffeomorphisms – expressly forbids a priori individuation of the points of a spacetime manifold as spatio-temporal events.

Although unfairly under-appreciated (especially within approaches to 'quantum gravity' through 'string' or 'M' type theories, practically all of which being guilty of presupposing one form or another of blatantly unjustified non-dynamical background structure (Rovelli 1997, Banks 1998a, b, Polchinski 1998, Sen 1998, Smolin 1998)), this is one of the most fundamental metaphysical tenets of general relativity. In this respect, contrary to what is often asserted following Kretschmann (1917), the principle of general covariance is far from being physically vacuous. For instance, the potency of general covariance is strikingly manifest in the following circumstance: if A_{μ} is a vector field on a general relativistic manifold \mathcal{M} , then, unlike the situation in electromagnetism discussed above, the value $A_{\mu}(x)$ at a particular point $x \in \mathcal{M}$ has no invariant physical meaning. This is because the point x can be actively transposed around by the action of the diffeomorphism group Diff(\mathcal{M}), robbing it of any individuality of its own.

Of course, individuation of spacetime points can be achieved by a fixed 'gauge choice' – that is to say, by specification of a particular metric tensor field $g^{\mu\nu}$ out of the entire equivalence class of fields $\{g^{\mu\nu}\}$ related by gauge transformations, but that would be at odds with general covariance. In fact, if there are any non-dynamical structures present, such as the globally specified Minkowski metric tensor field $\eta^{\mu\nu}$ of special relativity, then the impact of general covariance is severely mitigated. This is because the non-dynamical Minkowski metric tensor field, for example, can be used to introduce a family of global inertial coordinate systems (or 'inertial individuating fields' (Stachel 1993)) that can be transformed into each other by the (extended) Poincaré group of isometries of the metric: $\pounds_{\mathbf{x}}\eta^{\mu\nu}=0$, where $\pounds_{\mathbf{x}}$ denotes the Lie derivative, with the Killing vector field \mathbf{x}^{α} being a generator of the Poincaré group of transformations (Wald 1984). These inertial coordinates in turn can be used to set apart a point q from a point p of a manifold, bestowing apriori spatio-temporal individuality to the points of the manifold (Wald 1984). For this reason, Stachel (1993) and Wald (1984), among others, strengthen the statement of general covariance by a condition – explicitly added to the usual requirement of tensorial form for the law-like equations of physics – that there should not be any preferred individuating fields in spacetime other than, or independent of, the dynamically determined metric tensor field $g^{\mu\nu}$. Here preferred or background fields are understood to be the ones which affect the dynamical objects of a theory, but without being affected by them in return. They thereby provide non-dynamical backdrops for the dynamical processes. I shall return below to this issue of the background structure in spacetime.

In the light of this discussion, and in response to the lack of consensus on the meaning of general covariance in the literature (Norton 1993), let me end this section by proposing a litmus test for general covariance – formulated at the level of theory as a whole – which captures its true physical and metaphysical essence.

A litmus test for general covariance: A given theory may qualify to be called generally covariant if and only if the points of the spacetime 4-manifold, or a more general N-manifold, belonging to any model of the theory do not possess physically meaningful a priori individuality of their own.

(A model of a theory is a set of dynamical variables constituting a particular solution to the dynamical equations of the theory, and may, in general, also contain non-dynamical structures.) Admittedly, this is not a very practical elucidation of the principle, but it does exclude theories which are not truly generally covariant in the sense discussed above. In particular, it excludes *all* of the 'string' or 'M' type theories known to date, since they all presuppose individuation-condoning background structure of one form or another. (For a recent attempt to overcome this potentially detrimental deficiency of M-theory, see (Smolin 1998).)

3. Inadequacies of the Orthodox Quantum Theory of Measurement:

Even if we tentatively ignore the issue of individuation of spatio-temporal events, there exist a further concern regarding the notion of definite events in the quantum domain. In quantum theories (barring a few approaches to 'quantum gravity') one usually takes spacetime to be a fixed continuum whose constituents are the 'events' at points of space at instants of time. What is implicit in this assumption is the classicality or definiteness of the events. However, according to quantum mechanics, in general the notions such as 'here' and 'now' could have only indefinite or potential meaning. Further, if the conventional quantum framework is interpreted as universally applicable, objective (i.e., non-anthropocentric), and complete (Einstein et al. 1935), then, as pointed out above in the Introduction, the linear nature of quantum dynamics gives rise to some serious conceptual difficulties collectively known as 'the measurement problem' (Schrödinger 1935, Shimony 1963, Wheeler and Zurek 1983, Bell 1990). These difficulties make the notion of definite or actual events in the quantum world quite problematic, if not entirely meaningless (Jauch 1968, Haag 1990, 1992, Shimony 1993b). In particular, they render the orthodox quantum theory of measurement inadequate to explain the prolific occurrences of actual events in the 'macroscopic' domain, such as the sparks in a scintillation counter.

To elucidate the measurement problem, let us consider a highly schematized 'ideal measurement' type situation. Let Σ^s and Σ^A be two quantum mechanical systems constituting a *closed* composite system $\Sigma = \Sigma^s + \Sigma^A$ with their physical states represented by the rays corresponding to normalized vectors in the Hilbert spaces \mathcal{H}^s , \mathcal{H}^A , and $\mathcal{H}^{\Sigma} = \mathcal{H}^s \otimes \mathcal{H}^A$, respectively. Suppose now one wants to obtain the value of a dynamical variable corresponding to some property of the system Σ^s by means of the system Σ^A , which serves as a measuring apparatus. If this dynamical variable of Σ^s is represented by a self-adjoint operator Ω^s in the Hilbert space \mathcal{H}^s with the eigenvalue equation

$$\Omega^{S} |\psi_{i}\rangle = \omega_{i} |\psi_{i}\rangle \tag{3.1}$$

for some basis $\{|\psi_j\rangle\}\in\mathcal{H}^S$ ($\omega_i\neq\omega_j$ if $i\neq j$), then, for Σ^A to serve the purpose of measuring the value of a property of Σ^S in a state $|\psi_j\rangle$, there must be a vector $|\varphi_o\rangle$ in \mathcal{H}^A representing the ground state of the system Σ^A such that after a quantum mechanical interaction between Σ^S and Σ^A the resulting final state of the composite system $\Sigma = \Sigma^S + \Sigma^A$ is of the form $|\psi_j\rangle \otimes |\varphi_j\rangle$, where $\{|\varphi_j\rangle\}\in\mathcal{H}^A$ represents the set of indicator eigenstates

$$Q^{A} |\varphi_{j}\rangle = q_{j} |\varphi_{j}\rangle \tag{3.2}$$

 $(q_i \neq q_j \text{ if } i \neq j)$ constituting a basis in \mathcal{H}^A with $Q^A \in \mathcal{H}^A$ as a self-adjoint operator corresponding to a dynamical variable representing the 'indicator' property of the apparatus system Σ^A . Once such a correlation between the two subsystems is established, one can unequivocally infer the value ω_j corresponding to the property in question of the system Σ^S in the state $|\psi_j\rangle$ from the eigenvalue q_j of the indicator variable Q^A . For the sake of simplicity I have assumed here that Σ^A can only be in one of a discrete, non-degenerate set

of eigenstates $\{|\varphi_j\rangle\}$ of the indicator variable Q^A , and that the measurement interaction is of an 'ideal' type – i.e., the one which preserves the identity of the system Σ^S of interest as well as that of the system Σ^A treated as the measuring instrument.

So far the procedure described to infer the value of some property of a given quantum system does not involve any ambiguity. Unfortunately, this is not the case for more realistic initial states of the system of interest. In general, the initial state of the system Σ^s will not be an eigenstate but a superposition state of the form

$$|\psi\rangle = \sum_{j=1}^{N} \lambda_j |\psi_j\rangle,$$
 (3.3)

where the scalar coefficients $\lambda_j \in \mathbb{C}$, more than one being non-zero, satisfy $\sum_{j=1}^N |\lambda_j|^2 = 1$ with $N \equiv \dim \mathcal{H}^s$. The post-interaction entangled state of the composite system

$$|\Psi\rangle = \sum_{j=1}^{N} \lambda_j |\psi_j\rangle \otimes |\varphi_j\rangle$$
 (3.4)

dictated by the linear nature of quantum dynamics can now be seen to have generated an anomaly defying the very purpose of measurement. For, the resultant state (3.4) now itself is a superposition of N vectors $|\psi_j\rangle\otimes|\varphi_j\rangle$, each corresponding to a state in which the 'indicator' property of the apparatus system Σ^A has a different value q_j . In other words, the state $|\Psi\rangle$ as expressed in (3.4) implies that the 'indicator' property of the apparatus system is *indefinite* (a pointer on the dial of a detector does not point in any definite direction) obscuring the understanding of the observed actual occurrence of events such as a formation of a droplet in a cloud chamber, or a blackening of a silver grain on a photographic plate. The absurdity of this direct consequence of the linearity of quantum dynamics is well dramatized by Schrödinger in his gedanken-experiment involving a poor cat (1935), which ends up in a limbo between definite states of being alive and being dead.

In the conventional quantum mechanics this blatant contradiction with the apparent phenomenological facts about the occurrence of actual events is evaded by invoking an ad hoc postulate – the Projection Postulate, which in its simplest form is usually attributed to von Neumann. According to von Neumann's theory of measurement (1955), what has been described so far constitutes only the first stage of measurement. The second stage involves an instantaneous, discontinuous and acausal change, which cannot be described by the usual linear and reversible quantum dynamics, and is assumed to be accomplished by the development of the pure state $W \equiv \mathbb{P}_{|\Psi\rangle}$, $|\Psi\rangle = \sum_j \lambda_j |\psi_j\rangle \otimes |\varphi_j\rangle$, into the following proper mixture

$$W \longrightarrow W_{\text{or}} = \sum_{j} |\lambda_{j}|^{2} \mathbb{P}_{|\psi_{j}\rangle \otimes |\varphi_{j}\rangle} ,$$
 (3.5)

where $W, W_{\text{or}} \in \mathcal{T}(\mathcal{H}^{\Sigma})_{1}^{+}$ are positive normalized trace class operators in \mathcal{H}^{Σ} , and $\mathbb{P}_{|\eta\rangle}$ denotes an orthonormal projector onto the one-dimensional space spanned by a vector $|\eta\rangle$ in an appropriate Hilbert space. Note that now the state (3.5) of the composite system is consistent with, but of course does not imply, the phenomenological Born rule: If a quantum system Σ^{S} in an initial pure state represented by a unit vector $|\psi\rangle = \sum_{j} \lambda_{j} |\psi_{j}\rangle$ is measured by another quantum system Σ^{A} with indicator states $|\varphi_{j}\rangle$, then after the measurement interaction the state of the composite system $\Sigma = \Sigma^{S} + \Sigma^{A}$ is left in one of the pure states $|\psi_{k}\rangle \otimes |\varphi_{k}\rangle$ with probability $|\lambda_{k}|^{2}$, where the indicator state $|\varphi_{k}\rangle$ is correlated by the interaction with the eigenstate $|\psi_{k}\rangle$ of the dynamical variable Ω^{S} corresponding to some property of the system Σ^{S} .

A required third and the final stage in von Neumann's scheme of measurement involves the contentious 'ignorance interpretation of mixtures' (Beltrametti and Cassinelli 1981, Busch et al. 1991); for it is not yet clear how the measuring instrument Σ^A comes to exhibit a definite outcome – i.e., how the actual occurrence of a single definite event with corresponding relative frequency takes place out of the compendium of various possible events encoded in the mixture $W_{\rm or}$. This is because the nonlinear and stochastic transition $|\Psi\rangle \to |\psi_k\rangle \otimes |\varphi_k\rangle$ required by the Born rule implies the projection map $W \to W_{\rm or}$, but not vice versa, since

the mixture $W_{\rm or}$ in general does not have a unique decomposition in terms of the projector $\mathbb{P}_{|\psi_j\rangle\otimes|\varphi_j\rangle}$. In general, mixtures of numerous other projection operators may also be represented by the same statistical operator, making it impossible to say which 'basic' set of states $W_{\rm or}$ is a mixture of – a difficulty sometimes referred to as the 'preferred basis problem'. To put the rationale of difficulty in C^* -algebraic terms (Primas 1983), the space of resultant statistical states $W_{\rm or}$ does not form a 'simplex' in general (Choquet and Meyer 1963), disallowing the vectors $|\psi_j\rangle\otimes|\varphi_j\rangle$ to form a 'disjoint' set of vectors, and hence an ignorance or epistemic interpretation of such mixtures is not tenable.

The upshot clearly is that von Neumann's Projection Postulate is only a necessary but not sufficient condition for an unequivocal understanding of the occurrence of definite events. Even if we accept this *ad hoc* postulate unreservedly, the process of specific actualization out of the compendium of quantum mechanical potentialities remains completely obscure. Consequently, what is desperately needed is an unequivocal *physical* understanding underlying the non-unitary transition

$$\sum_{j=1}^{N} \lambda_j |\psi_j\rangle \otimes |\varphi_j\rangle \longrightarrow |\psi_k\rangle \otimes |\varphi_k\rangle. \tag{3.6}$$

As discussed in the Introduction above, despite a multitude of attempts with varied sophistication and predilections, no universally acceptable explanation – physical or otherwise – of this mysterious transition is as yet in sight. In the next section we shall see that Penrose's scheme provides precisely the much desired physical explanation for the transition, and compellingly so.

4. Penrose's Mechanism for the Objective State Reduction:

4.1. Motivation via a concrete example:

To illustrate Penrose's proposal within a concrete scenario, let us apply the above description of measurement procedure to a model interaction, within the nonrelativistic domain, in a specific representation – the coordinate representation. Let us begin by assuming a global inertial coordinate system whose origin is affixed at the center of the earth. Using this coordinate system, let the indicator variable Q^A represent the location q of the system Σ^A , which, say, has mass M, and let the dynamical variable Ω^S be a time-independent function of coordinate x and its conjugate momentum $-i\hbar\frac{\partial}{\partial x}$ of the system Σ^S exclusively. Further, let the mass M (i.e., the apparatus system Σ^A) be localized initially $(t < t_a)$ at q_1 , and let the measurement, which is to be achieved by moving the mass from q_1 to some other location, consist in the fact that if the value of $\Omega^S(x, -i\hbar\frac{\partial}{\partial x})$ is ω_1 then the location of the mass remains unchanged at q_1 whereas if it is $\omega_{j\neq 1}$ the mass is displaced from q_1 to a new location $q_{j\neq 1} \equiv q_1 + \omega_{j\neq 1}$. An interaction Hamiltonian which precisely accounts for such a process according to the conventional Schrödinger equation is (von Neumann 1955, d'Espagnat 1976)

$$H_{int}(t) = \beta(t) \widetilde{\Omega}^{S} P^{A}, \qquad (4.1)$$

where $\beta(t)$ is a smooth function of time $t \in \mathbb{R}$ with compact support $[t_a, t_b]$ satisfying

$$\int_{t_a}^{t_b} \beta(t) dt = 1, \qquad (4.2)$$

$$\widetilde{\Omega}^{S}(x, -i\hbar \frac{\partial}{\partial x}) \psi_{i} = \{\omega_{i} - \omega_{1} \delta(\omega_{1} - \omega_{i})\} \psi_{i} := \widetilde{\omega}_{i} \psi_{i}, \qquad (4.3)$$

and $P^A=-i\hbar \frac{\partial}{\partial q}$ with q being the indicator coordinate. The time-dependent Schrödinger equation

$$i\hbar \frac{\partial}{\partial t} |\Psi(t)\rangle = H(t) |\Psi(t)\rangle$$
 (4.4)

with the interaction Hamiltonian (4.1) has a general solution

$$\Psi(t) = \psi_i \, \varphi_i \, (q - q_1 - \alpha \, (t) \, \widetilde{\omega}_i) \, , \tag{4.5}$$

where φ_j is an arbitrary function of its argument determinable by initial conditions, and

$$\alpha(t) = \int_{0}^{t} \beta(t') dt' = \begin{cases} 0, & \forall t < t_a; \\ 1, & \forall t > t_b. \end{cases}$$

$$(4.6)$$

For the mass M localized at q_1 for $t < t_a$, the wave-function of the composite system $\Sigma = \Sigma^s + \Sigma^A$ is the product function

$$\Psi\left(t < t_{a}\right) = \psi_{i} \,\delta\left(q - q_{1}\right) \,, \tag{4.7}$$

and the function φ_j is identical to the delta-function. As a result, the state after $t = t_b$ is

$$\Psi(t > t_b) = \psi_i \,\delta(q - q_1 - \widetilde{\omega}_i) \tag{4.8}$$

according to Eq. (4.5). The interaction therefore induces a transition

$$\psi_j \, \delta \left(q - q_1 \right) \stackrel{H_{int}}{\longrightarrow} \psi_j \, \delta \left(q - q_1 - \widetilde{\omega}_j \right) \, . \tag{4.9}$$

In other words, using the definition (4.3) of $\widetilde{\omega}_i$,

$$\psi_{1} \, \delta \left(q - q_{1} \right) \quad \xrightarrow{H_{int}} \quad \psi_{1} \, \delta \left(q - q_{1} \right) \qquad \text{(location of } M \text{ unchanged)}$$
but
$$\psi_{j \neq 1} \, \delta \left(q - q_{1} \right) \quad \xrightarrow{H_{int}} \quad \psi_{j \neq 1} \, \delta \left(q - q_{j \neq 1} \right) \qquad \text{(location of } M \text{ shifted)} \,, \tag{4.10}$$

where recall that $q_{j\neq 1} \equiv q_1 + \omega_{j\neq 1}$. More generally, if the initial state of the quantum system Σ^s is a superposition state represented by

$$\sum_{j=1}^{N} \lambda_j \, \psi_j \, , \qquad \sum_{j=1}^{N} |\lambda_j|^2 = 1 \, , \tag{4.11}$$

then we have the above discussed Schrödinger's Cat (1935) type entanglement exhibiting superposition of the location-states of the mass at various positions:

$$\left[\sum_{j=1}^{N} \lambda_{j} \psi_{j}\right] \delta\left(q - q_{1}\right) \xrightarrow{H_{int}} \sum_{j=1}^{N} \lambda_{j} \psi_{j} \delta\left(q - q_{j}\right) \equiv \sum_{j=1}^{N} \lambda_{j} \psi_{j} \varphi_{j}. \tag{4.12}$$

In particular, if initially we have

$$\Psi(t < t_a) = (\lambda_2 \psi_2 + \lambda_3 \psi_3) \ \delta(q - q_1) , \quad |\lambda_2|^2 + |\lambda_3|^2 = 1 , \qquad (4.13)$$

then, after the impulsive interaction,

$$\Psi(t > t_b) = \lambda_2 \psi_2 \, \delta(q - q_2) + \lambda_2 \psi_2 \, \delta(q - q_2) , \qquad (4.14)$$

and the location of the mass will be indefinite between the two positions q_2 and q_3 . This of course is a perfectly respectable quantum mechanical state for the mass M to be in, unless it is a 'macroscopic' object and the two locations are macroscopically distinct. In that case the indefiniteness in the location of the mass dictated by the linearity of quantum dynamics stands in a blatant contradiction with the evident phenomenology of such objects.

4.2. The raison d'être of state reduction:

Recognizing this contradiction, Penrose, among others, has tirelessly argued that gravitation must be directly responsible for an *objective* resolution of this fundamental anomaly of quantum theory (1979, 1981, 1984, 1986, 1987, 1989, 367-371, 1993, 1994a, 1994b, 339-346, 1996, 1997, 1998). He contends that, since the self-gravity of the mass must also participate in such superpositions, what is actually involved here, in accordance with the principles of Einstein's theory of gravity, is a superposition of two entirely *different* spacetime geometries; and, when the two geometries are sufficiently different from each other, the unitary quantum mechanical description of the situation – i.e., the linear superposition of a 'macroscopic' mass prescribed by Eq. (4.14) – must breakdown² (or, rather, 'decay'), allowing nature to choose between one or the other of the two geometries.

To understand this claim, let me surface some of the hidden assumptions regarding spacetime structure underlying the time-evolution dictated by Eq. (4.4), which brought us to the state (4.14) in question. Recall that I began this section with an assumption of a globally specified inertial frame of reference affixed at the center of earth. Actually, this is a bit too strong an assumption. Since the Schrödinger equation (4.4) is invariant under Galilean transformations, all one needs is a *family* of such global inertial frames, each member of which is related to another by a Galilean transformation

$$t \longrightarrow t' = t + constant,$$

 $x^a \longrightarrow x'^a = O^a_b x^b + v^a t + constant, \qquad (a,b=1,2,3),$ (4.15)

where $O^a_{\ b} \in SO(3)$ is a time-independent orthonormal rotation matrix (with Einstein's summation convention for like indices), and $\mathbf{v} \in \mathbb{R}^3$ is a time-independent spatial velocity. Now, as discussed at the end of section 2 above, existence of a global inertial frame grants a priori individuality to spacetime points – a point p_1 of a spacetime manifold can be set apart from a point p_2 using such inertial coordinates (Wald 1984, 6). Consequently, in the present scenario the concepts 'here' and 'now' have a priori meaning, and they can be taken as a part of any physical question (cf. section 2). In particular, it is meaningful to take location q_1 of the mass M to be a part of the initial state (4.13), since it can be set apart from any other location, such as the location q_2 or q_3 in the final state (4.14). If individuation of spatio-temporal events was not possible, then of course all of the locations, q_1 , q_2 , q_3 , etc., would have been identified with each other as one and the same location, and it would not have been meaningful to take q_1 as a distinct initial location of the mass (as elaborated in section 2 above, such an identification of all spacetime points is indeed what general covariance demands in full general relativity). Now, continuing to ignore gravity for the moment, but anticipating Penrose's reasoning when gravity is included, let us pretend, for the sake of argument, that the two components of the superposition in Eq. (4.14) correspond to two different (flat) spacetime geometries. Accordingly, let us take two separate inertial coordinate systems, one for each spacetime but related by the transformation (4.15), for separately describing the evolution of each of the two components of the superposition, with the initial location of the mass M being q_1 as prescribed in Eq. (4.13) – i.e., assume for the moment that each component of the superposition is evolving on its own, as it were, under the Schrödinger equation (4.4). Then, for the final superposed state (4.14) to be meaningful, a crucially important question would be: are these two time-evolutions corresponding to the two different spacetimes compatible with each other? In particular: is the time-translation operator $(\frac{\partial}{\partial t})$ in Eq. (4.4) the same for the two superposed evolutions – one displacing the mass M from q_1 to q_2 and the other displacing it from q_1 to q_3 ? Unless the two time-translation operators in the two coordinate systems are equivalent in some sense, we do not have a meaningful quantum gestation of the superposition (4.14). Now, since we are in the Galilean-relativistic domain, the two inertial frames assigned to the two spacetimes must be related by the transformation (4.15),

It is worth emphasising here that, as far as I can infer from his writings, Penrose is not committed to any of the existing proposals of nonlinear (e.g., Weinberg 1989) and/or stochastic (e.g., Pearle 1993) modifications of quantum dynamics (neither am I for that matter). Such proposals have their own technical and/or interpretational problems, and are far from being completely satisfactory. As discussed in the Introduction, Penrose's proposal, by contrast, is truly minimalist. Rather than prematurely proposing a theory of quantum state reduction, he simply puts forward a rationale why his heuristic scheme for the actualization potentialities must inevitably be a builtin feature of the sought-for 'final theory'.

which, upon using the chain rule (and setting $O \equiv 1$ for simplicity), yields

$$\frac{\partial}{\partial x'^a} = \frac{\partial}{\partial x^a}, \quad \text{but} \quad \frac{\partial}{\partial t'} = \frac{\partial}{\partial t} - v^a \frac{\partial}{\partial x^a}.$$
 (4.16)

Thus, the time-translation operators are *not* the same for the two spacetimes (cf. Penrose 1996, 592-593). As a result, in general, unless ${\bf v}$ identically vanishes everywhere, the two superposed time-evolutions are *not* compatible with each other (see subsection 5.2, however, for a more careful analysis). The difficulty arises for the following reason. Although in this Galilean-relativistic domain the individuality of spacetime points in a given spacetime is rather easy to achieve, when it comes to two entirely *different* spacetimes there still remains an ambiguity in registering the fact that the location, say q_2 , of the mass in one spacetime is 'distinct' from its location, say q_3 , in the other spacetime. On the other hand, the location q_2 must be unequivocally distinguishable from the location q_3 for the notion of superposition of the kind (4.14) to have any unambiguous physical meaning. Now, in order to meaningfully set apart a location q_2 in *one* spacetime from a location q_3 in *another*, a point-by-point identification of the two spacetimes is clearly necessary. But such a *pointwise* identification is quite ambiguous for the two spacetimes under consideration, as can be readily seen form Eq. (4.15), unless the arbitrarily chosen relative spatial velocity ${\bf v}$ is set to identically vanish everywhere (i.e., not just locally). Of course, in the present scenario, since we have ignored gravity, nothing prevents us from setting ${\bf v} \equiv 0$ everywhere – i.e., by simply taking nonrotating coordinate systems with constant spatial distance between them – and the apparent difficulty completely disappears, yielding

$$\frac{\partial}{\partial t'} \equiv \frac{\partial}{\partial t} \,. \tag{4.17}$$

Therefore, as long as gravity is ignored, there is nothing wrong with the quantum mechanical time-evolution leading to the superposed state (4.14) from the initial state (4.13), since all of the hidden assumptions exposed in this paragraph are more than justified.

The situation becomes dramatically obscure, however, when one attempts to incorporate gravity in the above scenario in full accordance with the principles of general relativity³. To appreciate the central difficulty, let us try to parallel considerations of the previous paragraph with due respect to the ubiquitous general-relativistic features of spacetime⁴. To begin with, once gravity is included, even the initial state (4.13) becomes meaningless because any location such as q_1 loses its a priori meaning. Recall from section 2 that in general relativity, since neither global topological structure of spacetime nor local individuality of spatio-temporal events has any meaning until a specific metric tensor field is dynamically determined, the concepts 'here' and 'now' can only be part of the answer to a physical question. On the other hand, the initial state (4.13) specifying the initial location q_1 of the mass M is part of the question itself regarding the evolution of the mass. Thus, from the general-relativistic viewpoint – which clearly is the correct viewpoint for a 'large enough' mass – the statement (4.13) is entirely meaningless. In practice, however, for the nonrelativistic situation under consideration, the much more massive earth comes to rescue, since it can be used to serve as an external frame of reference providing prior – albeit approximate – individuation of spacetime points (Rovelli 1991). For the sake of argument, let us be content with such an approximate specification of the initial location q_1 of the mass, and ask: what role would the general-relativistic features of spacetime play in the evolution of this mass either from q_1 to q_2 or from q_1 to q_3 , when these two evolutions are viewed separately - i.e., purely 'classically'? Now, since the self-gravity of the mass must also be taken into account here, and since each of the two evolutions would incorporate the self-gravitational effects in its own distinct manner to determine its own overall a posteriori spacetime geometry in accordance with the

³ It is worth noting here that the conventional 'quantum gravity' treatments are of no help in the conceptual issues under consideration. Indeed, as Penrose points out (1996, 589), the conventional attitude is to treat superpositions of different spacetimes in merely formal fashion, in terms of complex functions on the space of 3- or 4-geometries, with no pretence at conceptual investigation of the physics that takes place within such a formal superposition.

⁴ Within our nonrelativistic domain, a more appropriate spacetime framework is of course that of Newton-Cartan theory (Christian 1997). This framework will be taken up in a later more specialized discussion, but for now, for conceptual clarity, I rather not deviate from the subtleties of the full general-relativistic picture of spacetime.

dynamical principles of Einstein's theory, to a good degree of classical approximation there will be essentially two distinct spacetime geometries associated with these two evolutions. Actually, as in the case of initial location q_1 , the two final locations, q_2 and q_3 , would also acquire physical meaning only a posteriori via the two resulting metric tensor fields – say $g_2^{\mu\nu}$ and $g_3^{\mu\nu}$, respectively, since the individuation of the points of each of these two spacetimes becomes meaningful only a posteriori by means of these metric tensor fields. It is of paramount importance here to note that, in general, the metric tensor fields $g_{\alpha}^{\mu\nu}$ and $g_{\alpha}^{\mu\nu}$ would represent two strictly separate spacetimes with their own distinct global topological and local causal structures. To dramatize this fact by means of a rather extreme example, note that one of the two components of the superposition leading to Eq. (4.14) might, in principle, end up having evolved into something like a highly singular Kerr-Newman spacetime, whereas the other one might end up having evolved into something like a non-singular Robertson-Walker spacetime. This observation is crucial to Penrose's argument because, as we did in the previous paragraph for the non-gravitational case, we must now ask whether it is meaningful to set apart one location of the mass, say q_2 , from another, say q_3 , in order for a superposition such as (4.14) to have any unambiguous physical meaning. And as before, we immediately see that in order to be able to distinguish the two locations of the mass – i.e., to register the fact that the mass has actually been displaced from the initial location q_1 to a final location, say q_2 , and not to any other location, say q_3 – a point-by-point identification of the two spacetimes is essential. However, in the present general-relativistic picture such a pointwise identification is utterly meaningless, especially when the two geometries under consideration are 'significantly' different from each other. As a direct consequence of the principle of general covariance, there is simply no meaningful way to make a pointwise identification between two such distinct spacetimes in general relativity. Since the theory makes no a priori assumption as to what the spacetime manifold is and allows the Lorentzian metric tensor field to be any solution of Einstein's field equations, the entire causal structure associated with a general-relativistic spacetime is dynamical and not predetermined (cf. section 2). In other words, unlike in special relativity and the case considered in the previous paragraph, there is simply no isometry group underlying the structure of general relativity which could allow existence of a preferred family of inertial reference frames that may be used, first, to individuate the points of each spacetime, and then to identify one spacetime with another point-by-point. Furthermore, the lack of an isometry group means that, in general, there are simply no Killing vector fields of any kind in a general-relativistic spacetime, let alone a time-like Killing vector field analogous to the time-translation operator ' $\frac{\partial}{\partial t}$ ' of the non-gravitational case considered above (cf. Eq. (4.16)). Therefore, in order to continue our argument, we have to make a further assumption: We have to assume, at least, that the spacetimes under consideration are actually two reasonably well-defined 'stationary' spacetimes with two time-like Killing vector fields corresponding to the time-symmetries of the two metric tensor fields $g_2^{\mu\nu}$ and $g_3^{\mu\nu}$, respectively. These Killing vector fields, we hope, would generate time-translations needed to describe the time-evolution analogous to the one provided by the operator $(\frac{\partial}{\partial t})$ in the non-gravitational case. However, even this drastic assumption hardly puts an end to the difficulties involved in the notion of time-evolution leading to a superposition such as (4.14). One immediate difficulty is that these two Killing vector fields generating the time-evolution are completely different for the two components of the superposition under consideration. Since they correspond to the time-symmetries of two essentially distinct spacetimes, they could hardly be the same. As a result, the two Killing vector fields represent two completely different causal structures, and hence, if we insist on implementing them, the final state corresponding to Eq. (4.14) would involve some oxymoronic notion such as 'superposition of two distinct causalities'. Incidentally, this problem notoriously reappears in different guises in various approaches to 'quantum gravity', and it is sometimes referred to as the 'problem of time' (Kuchař 1991, 1992, Isham 1993, Belot and Earman 1999). In summary, for a 'large enough' mass M, the final superposed state such as (4.14) is fundamentally and hopelessly meaningless.

4.3. Phenomenology of the objective state reduction:

In the previous two paragraphs we saw two extreme cases. In the first of the two paragraphs we saw that, as long as gravity is ignored, the notion of quantum superposition is quite unambiguous, thanks to the availability of a priori and exact pointwise identification between the two 'spacetimes' into which a mass M could evolve. However, since the ubiquitous gravitational effects cannot be ignored for a 'large enough' mass, in the last paragraph we saw that a notion of superposition within two general-relativistic spacetimes is completely meaningless. Thus, a priori and exact pointwise identification of distinct spacetimes – although

expressly forbidden by the principle of general covariance – turns out to be an essential prerequisite for the notion of superposition. In other words, the superposition principle is not as fundamental a principle as the adherents of orthodox quantum mechanics would have us believe; it makes sense only when the other most important principle – the principle of general covariance – is severely mitigated. By contrast, of course, as a result of formidable difficulties encountered in attempts to construct a $Diff(\mathcal{M})$ -invariant quantum field theory (Rovelli 1997), it is not so unpopular to adhere that *active* general covariance may be truly meaningful only at the classical general-relativistic level – i.e., when the superposition principle is practically neutralized.

To bridge this gulf between our two most basic principles at least phenomenologically, Penrose invites us to contemplate an intermediate physical situation for which the notion of quantum superposition is at best approximately meaningful. In a nutshell, his strategy is to first consider a 'superposition' such as (4.14) with gravity included, but nevertheless a priori pointwise identification of the two spacetimes corresponding to the two components of the superposition retaining some approximate meaning, and then, after putting a practical measure on this approximation, use this measure to obtain a heuristic formula for the collapse time of this superposition. Here is how this works: Consider two well-defined quantum states represented by $|\Psi_2\rangle$ and $|\Psi_3\rangle$ (analogous to the states (4.8)), each stationary on its own and possessing the same energy E:

$$i\hbar \frac{\partial}{\partial t} |\Psi_2\rangle = E |\Psi_2\rangle , \qquad i\hbar \frac{\partial}{\partial t} |\Psi_3\rangle = E |\Psi_3\rangle .$$
 (4.18)

In standard quantum mechanics, when gravitational effects are ignored, linearity dictates that any superposition of these two stationary states such as

$$|\mathcal{X}\rangle = \lambda_2 |\Psi_2\rangle + \lambda_3 |\Psi_3\rangle \tag{4.19}$$

(cf. Eq. (4.14)) must also be stationary, with the same energy E:

$$i\hbar \frac{\partial}{\partial t} |\mathcal{X}\rangle = E |\mathcal{X}\rangle.$$
 (4.20)

Thus, quantum linearity necessitates a complete degeneracy of energy for superpositions of the two original states. However, when the gravitational fields of two different mass distributions are incorporated in the representations $|\Psi_2\rangle$ and $|\Psi_3\rangle$ of these states, a crucial question arises: will the state $|\mathcal{X}\rangle$ still remain stationary with energy E? Of course, when gravity is taken into account, each of the two component states would correspond to two entirely different spacetimes with a good degree of classical approximation, whether or not we assume that they are reasonably well-defined stationary spacetimes. Consequently, as discussed above, the time-translation operators such as $(\frac{\partial}{\partial t})$ corresponding to the action of the time-like Killing vector fields of these two spacetimes would be completely different form each other in general. They could only be the same if there were an unequivocal pointwise correspondence between the two spacetimes. Let us assume, however, that these two Killing vector fields are not too different from each other for the physical situation under consideration. In that case, there would be a slight – but essential – ill-definedness in the action of the operator $(\frac{\partial}{\partial t})$ when it is employed to generate a superposed state such as (4.19), and this ill-definedness would be without doubt reflected in the energy E of this state. One can use this ill-definedness in energy, ΔE , as a measure of instability of the state (4.19), and postulate the life-time of such a 'stationary' superposition – analogous to the half-life of an unstable particle – to be

$$\tau = \frac{\hbar}{\Delta E} \,, \tag{4.21}$$

with two decay modes being the individual states $|\Psi_2\rangle$ and $|\Psi_3\rangle$ with relative probabilities $|\lambda_2|^2$ and $|\lambda_3|^2$, respectively. Clearly, when there is an exact pointwise identification between the two spacetimes, $\Delta E \to 0$, and the collapse of the superposition never happens. On the other hand, when such an identification is ambiguous or impossible, inducing much larger ill-definedness in the energy, the collapse is almost instantaneous.

A noteworthy feature of the above formula is that it is independent of the speed of light c, implying that it remains valid even in the nonrelativistic domain (cf. Figure 1 above and (Penrose 1994, 339, 1996, 592)).

Further, in such a Newtonian approximation, the ill-definedness ΔE (for an essentially static situation) turns out to be proportional to the gravitational self-energy of the difference between the mass distributions belonging to the two components of the superposition (Penrose 1996). Remarkably, numerical estimates (Penrose 1994, 1996) based on such Newtonian models for life-times of superpositions turn out to be strikingly realistic. For instance, the life-time of superposition for a proton works out to be of the order of a few million years, whereas a water droplet – depending on its size – is expected to be able to maintain superposition only for a fraction of a second. Thus, the boundary near which the reduction time is of the order of seconds is precisely the phenomenological quantum-classical boundary of our corroborative experience⁵.

As I alluded to towards the end of section 3, an important issue in any quantum measurement theory is the 'preferred basis problem'. The difficulty is that, without some further criterion, one does not know which states from the general compendium of possibilities are to be regarded as the 'basic' (or 'stable' or 'stationary') states and which are to be regarded as essentially unstable 'superpositions of basic states' – the states which are to reduce into the basic ones. Penrose's suggestion is to regard – within Newtonian approximation – the stationary solutions of what he calls the Schrödinger-Newton equation as the basic states (Penrose 1998, Moroz et al. 1998, Tod and Moroz 1998). I shall elaborate on this equation (which I have independently studied in (Christian 1997)) in the next section.

4.4. A different measure of deviation from quantum mechanics:

As an aside, let me propose in this subsection a slightly different measure for the lack of exact pointwise identification between the two spacetimes under consideration. In close analogy with the above assumption of stationarity, let us assume that there exists a displacement isometry in each of the two spacetimes, embodied in the Killing vector fields \mathbf{x}_2 and \mathbf{x}_3 respectively – i.e., let $\pounds_{\mathbf{x}_2} g_2^{\ \mu\nu} = 0 = \pounds_{\mathbf{x}_3} g_3^{\ \mu\nu}$, where $\pounds_{\mathbf{x}}$ denotes the Lie derivative with Killing vector fields \mathbf{x}_2^{α} and \mathbf{x}_3^{α} as the generators of the displacement symmetry. Further, as before, let us assume that at least some approximate pointwise identification between these two spacetimes is meaningful. As a visual aid, one may think of two nearly congruent coordinate grids, one assigned to each spacetime. Then, \grave{a} la Penrose, I propose a measure of incongruence between these two spacetimes to be the dimensionless parameter $\mathbf{d}^{\sigma}\mathbf{d}_{\sigma}$, taking values between zero and unity, $0 \leq \mathbf{d}^{\sigma}\mathbf{d}_{\sigma} \leq 1$, with

$$d^{\sigma} := x_2^{\alpha} \nabla_{\alpha} x_3^{\sigma} - x_3^{\alpha} \nabla_{\alpha} x_2^{\sigma}. \tag{4.22}$$

As it stands, this quantity is mathematically ill-defined since the Killing vectors \mathbf{x}_2^{α} and \mathbf{x}_3^{α} describe the same displacement symmetry in two quite distinct spacetimes. However, if we reinterpret these two vectors as describing two slightly different symmetries in one and the same spacetime, then the vector field \mathbf{d}^{σ} is geometrically well-defined, and it is nothing but the commutator Killing vector field (Misner *et al.* 1973, 654) corresponding to the two linearly independent vectors \mathbf{x}_2^{α} and \mathbf{x}_3^{α} . In other words, \mathbf{d}^{σ} then is simply a measure of incongruence between the two coordinates adapted to simultaneously describe symmetries corresponding to both \mathbf{x}_2^{α} and \mathbf{x}_3^{α} within this single spacetime. This measure can now be used to postulate a gravity-induced deviation from the orthodox quantum commutation relation for the position and momentum of the mass M:

$$[Q, P] = i\hbar \left\{ 1 - d^{\sigma} d_{\sigma} \right\}. \tag{4.23}$$

Clearly, when there is an exact pointwise correspondence between the two spacetimes – i.e., when the Killing vector fields \mathbf{x}_{2}^{α} and \mathbf{x}_{3}^{α} are strictly identified and $\mathbf{d}^{\sigma}\mathbf{d}_{\sigma} \equiv 0$, we recover the standard quantum mechanical commutation relation between the position and momentum of the mass. On the other hand, when – for a 'large enough' mass – the quantity $\mathbf{d}^{\sigma}\mathbf{d}_{\sigma}$ reaches order unity, the mass exhibits essentially classical behaviour.

It should be noted that, independently of Penrose, Diósi has also proposed the same formula (4.21) for the collapse time (1989), but he arrives at it from a rather different direction. Penrose's scheme should also be contrasted (Penrose 1996) with the 'semi-classical approaches' to 'quantum gravity' (e.g., Kibble 1981), which are well-known to be inconsistent (Eppley and Hannah 1977, Wald 1984, 382-383, Anandan 1994). Recently, Anandan (1998) has generalized Penrose's Newtonian expression for ΔE to a similar expression for an arbitrary superposition of relativistic, but weak, gravitational fields, obtained in the gravitational analogue of the Coulomb gauge in a linearized approximation applied to the Lorentzian metric tensor field (cf. subsection 5.3 for further comments).

Thus, the parameter $d^{\sigma}d_{\sigma}$ provides a good measure of ill-definedness in the canonical commutation relation due to a Penrose-type incongruence, but now between the *displacement* symmetries of the two spacetimes.

4.5. Penrose's proposed experiment:

Finally, let me end this section by describing a variant of a realizable experiment proposed by Penrose to corroborate the contended 'macroscopic' breakdown of quantum mechanics (1998). The present version of the experiment due to Hardy (1998) is – arguably – somewhat simpler to perform. There are many practical problems in both Penrose's original proposal and Hardy's cleverer version of it (contamination due to the ubiquitous decoherence effects being the most intractable of all problems), but such practical problems will not concern us here (cf. Penrose 1998, Hardy 1998). Further, the use of a photon in the described experiment is for convenience only; in practice it may be replaced by any neutral particle, such as an ultracold atom of a suitable kind.

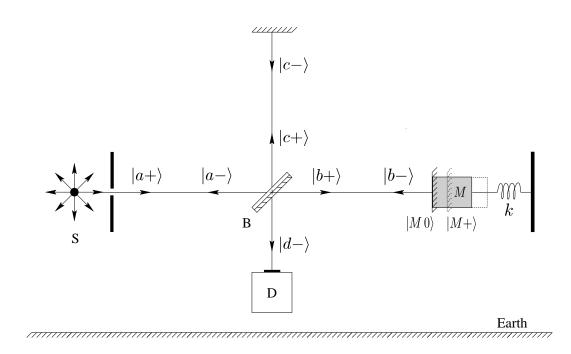


Figure 3: Hardy's version of Penrose's proposed experiment: In an interferometric arrangement, a beam-splitter, B, is placed in the 'path' of an incident photon emanating form a source S. A horizontally movable mass M is attached to the wall opposite to S by means of a restoring device with a spring constant k. There are two reflecting mirrors – one of them affixed on the mass and the other one at the end of the vertical arm of the interferometer, both being at an exactly equal distance from the beam splitter. The earth provides a frame of reference, and the final destination of interest for the photon is the detector D.

The basic experimental set-up is described in Figure 3. The system consists of two objects: a 'photon' and a 'macroscopic' object of mass M, which in Penrose's version is a small Mössbauer crystal with about 10^{15} nuclei. The objective of the arrangement is to render the 'macroscopic' mass in a superposition of two macroscopically distinct positions, as in the state (4.14) above. The '+' or '-' sign in the photon states (such as $|a+\rangle$ or $|c-\rangle$), respectively, indicate a forward or backward motion along a given 'path'. For now, we simply look at the arrangement in a purely orthodox, quantum mechanical fashion. Then, the following

transformations of the photon states due to a beam-splitter may be adopted from quantum optics:

$$|a\pm\rangle \longleftrightarrow \frac{1}{\sqrt{2}} \{ |b\pm\rangle + |c\pm\rangle \}$$

$$|d\pm\rangle \longleftrightarrow \frac{1}{\sqrt{2}} \{ |c\pm\rangle - |b\pm\rangle \}, \qquad (4.24)$$

with inverse relations being

$$|c\pm\rangle \longleftrightarrow \frac{1}{\sqrt{2}} \{ |a\pm\rangle + |d\pm\rangle \}$$

$$|b\pm\rangle \longleftrightarrow \frac{1}{\sqrt{2}} \{ |a\pm\rangle - |d\pm\rangle \}. \tag{4.25}$$

If the initial state of the incident photon is taken to be $|a+\rangle$, and the initial (or unmoved) state of the mass M is denoted by $|M0\rangle$, then the initial state of the *closed* composite system is the product state

$$|a+\rangle \otimes |M0\rangle$$
. (4.26)

As the photon passes through the beam-splitter, this composite initial state evolves into

$$\frac{1}{\sqrt{2}}\{|b+\rangle + |c+\rangle\} \otimes |M0\rangle. \tag{4.27}$$

Now, in the absence of the beam-splitter, if the photon happens to be in the horizontal 'path', then it would reflect off the mirror affixed on the mass, giving it a minute momentum in the '+' direction. On the other hand, if the photon is arranged to be in the vertical 'path', then it would simply reflect off the second mirror at the end of that path, without affecting the mass. The net result of these two alternatives in the presence of the beam-splitter, viewed quantum mechanically, is encoded in the state

$$\frac{1}{\sqrt{2}}\{|b-\rangle\otimes|M+\rangle + |c-\rangle\otimes|M0\rangle\}. \tag{4.28}$$

Since each of the two options in this superposition would lead the photon back towards the beam-splitter, the composite state (4.28) – as the photon passes again through the beam-splitter – will evolve into

$$\frac{1}{2} \Big[\{ |a-\rangle - |d-\rangle \} \otimes |M+\rangle + \{ |a-\rangle + |d-\rangle \} \otimes |M0\rangle \Big]. \tag{4.29}$$

Now, our goal here is to generate a Penrose-type superposition of the mass M. Therefore, at this stage we isolate only those sub-states for which the photon could be detected by the detector D. Thus selected from (4.29), we obtain

$$\frac{1}{\sqrt{2}}\{|M0\rangle - |M+\rangle\}\tag{4.30}$$

for the state of the mass, isolating it in the desired, spatially distinct, 'macroscopic' superposition. After some minute lapse of time, say Δt , the spring will bring the mass back to its original position with its momentum reversed, and thereby transform the above state into

$$\frac{1}{\sqrt{2}}\{|M0\rangle - |M-\rangle\},\tag{4.31}$$

where $|M-\rangle$ is the new state of M with its momentum in the '-' direction (not shown in the figure).

At this precise moment, in order to bring about decisive statistics, we send *another* photon from S into the interferometer which, upon passing through the beam-splitter, will produce the product state

$$\frac{1}{2} \left\{ |b+\rangle + |c+\rangle \right\} \otimes \left\{ |M0\rangle - |M-\rangle \right\}. \tag{4.32}$$

Just as before, the four terms of this state will now evolve on their own, and, after a recoil of the photon from the two mirrors, the composite state will become

$$\frac{1}{2} \left[|b-\rangle \otimes |M+\rangle + |c-\rangle \otimes |M0\rangle - |b-\rangle \otimes |M0\rangle - |c-\rangle \otimes |M-\rangle \right]. \tag{4.33}$$

It is crucial to note here that, in the third term, the momentum of the mass has been reduced to zero by the interaction so that both the second and third terms have the same state $|M0\rangle$ for the mass. Finally, the evolution of the photon back through the beam-splitter will render the composite system to be in the state

$$\frac{1}{2\sqrt{2}}\left\{\left|a-\right\rangle - \left|d-\right\rangle\right\} \otimes \left|M+\right\rangle + \frac{1}{\sqrt{2}}\left|d-\right\rangle \otimes \left|M0\right\rangle - \frac{1}{2\sqrt{2}}\left\{\left|a-\right\rangle + \left|d-\right\rangle\right\} \otimes \left|M-\right\rangle. \tag{4.34}$$

Thus, quantum mechanics predicts that the probability of detecting a photon in the detector D is 75%.

On the other hand, if the 'macroscopic' superposition of the mass such as (4.30) has undergone a Penrose-type process of state reduction, then the state of the mass just before the second photon is sent in would not be (4.31) but a proper mixture of $|M0\rangle$ and $|M-\rangle$. As a result, instead of (4.32), the overall disjoint state after the photon has passed through the beam-splitter would simply be

$$\frac{1}{\sqrt{2}}\{|b+\rangle + |c+\rangle\} \otimes |M0\rangle \quad \text{or} \quad \frac{1}{\sqrt{2}}\{|b+\rangle + |c+\rangle\} \otimes |M-\rangle, \tag{4.35}$$

without any quantum coherence between the two alternatives. As the photon is reflected off the two mirrors and passed again through the beam-splitter, these two 'classical' alternatives – instead of (4.34) – would evolve independently into the final disjoint state

$$\frac{1}{2}\Big[\big\{\,|a-\rangle-|d-\rangle\big\}\otimes|M+\rangle+\big\{\,|a-\rangle+|d-\rangle\big\}\otimes|M0\rangle\Big]\quad\text{or}\quad \frac{1}{2}\Big[\big\{\,|a-\rangle-|d-\rangle\big\}\otimes|M0\rangle+\big\{\,|a-\rangle+|d-\rangle\big\}\otimes|M-\rangle\Big]\,. \tag{4.36}$$

Consequently, if Penrose's proposal is on the right track, then, after the photon passes through the beam-splitter second time around, it would go to the detector only 50% of the time and not 75% of the time as quantum mechanics predicts. Practical difficulties aside (Penrose 1998, Hardy 1998), this is certainly a refutable proposition (especially because the commonly held belief concerning decoherence (Kay 1998) – i.e., a belief that a strong coupling to the environment inevitably destroys the observability of quantum effects between macroscopically distinct states – is quite misplaced, as emphasised by Leggett (1998)).

5. A Closer Look at Penrose's Proposal within Newton-Cartan Framework:

My main goal in this section is, first, to put forward a delicate argument that demonstrates why Penrose's experiment – as it stands – is not adequate to corroborate the signatures of his contended gravity-induced quantum state reduction, and then to briefly discuss a couple of decisive experiments which *would* be able to divulge the putative breakdown of quantum mechanics along the line of his reasoning.

5.1. An orthodox analysis within strictly Newtonian domain:

In order to set the stage for my argument, let us first ask whether one can provide an orthodox quantum mechanical analysis of the physics underlying Penrose's proposed experiment. As it turns out, one can indeed provide such an orthodox treatment. In this subsection I shall outline one such treatment, which will not only direct us towards pinpointing where and for what reasons Penrose's approach differs form the orthodox approach, but will also allow us to explore more decisive experiments compared to the one he has proposed.

Clearly, to respond to Penrose's overall conceptual scheme in orthodox mannar one would require a full-blown and consistent quantum theory of gravity, which, as we know, is not yet in sight (Rovelli 1998). If we concentrate, however, not on his overall conceptual scheme but simply on his proposed experiment, then we only require a nonrelativistic quantum theory of gravity (recall from the last section that the formula (4.21) does not depend on the speed of light). And, fortunately, such a theory does exist. Recently, I have been able to demonstrate (Christian 1997) that the covariantly described Newtonian gravity – the so-called Newton-Cartan gravity which duly respects Einstein's principle of equivalence – interacting with Galilean-relativistic matter (Schrödinger fields) exists as an exactly soluble system, both classically and quantum mechanically (cf. Figure 1). The significance of the resulting manifestly covariant unitary quantum field theory of gravity lies in the fact that it is the Newton-Cartan theory of gravity, and not the original Newton's theory of gravity, that is the true Galilean-relativistic limit form of Einstein's theory of gravity. In fact, an alternative, historically counterfactual but logically more appropriate, formulation of general relativity is simply Newton-Cartan theory of gravity 'plus' the light-cone structure of the special theory of relativity. Newton's original theory in such a 'generally-covariant' Newton-Cartan framework emerges in an adscititiously chosen local inertial frame (modulo a crucially important additional restriction on the curvature tensor, as we shall see).

To begin the analysis, let us first look at the classical Newton-Cartan theory (for further details and extensive references consult section II of (Christian 1997)). Cartan's spacetime reformulation of the classical Newtonian theory of gravity can be motivated in exact analogy with Einstein's theory of gravity. The analogy works because the universal equality of the inertial and the passive gravitational masses is independent of the relativization of time, and hence is equally valid at the Galilean-relativistic level. As a result, it is possible to parallel Einstein's theory and reconstrue the trajectories of (only) gravitationally affected particles as geodesics of a unique, 'non-flat' connection Γ satisfying

$$\frac{d^2x^i}{dt^2} + \Gamma_j^{\ i}_{\ k} \frac{dx^j}{dt} \frac{dx^k}{dt} = 0 \tag{5.1}$$

in a coordinate basis, such that

$$\Gamma^{\mu}_{\nu\lambda} \equiv \Gamma^{\nu}_{\nu\lambda} + \Theta^{\mu}_{\nu\lambda} := \Gamma^{\nu}_{\nu\lambda} + h^{\mu\alpha} \nabla^{\nu}_{\alpha} \Phi t_{\nu\lambda} , \qquad (5.2)$$

with Φ representing the Newtonian gravitational potential relative to the freely falling observer field v, $\Gamma^{\nu}_{\nu}{}^{\mu}_{\lambda}$ representing the coefficients of the corresponding 'flat' connection (i.e., one whose coefficients can be made to vanish in a suitably chosen linear coordinate system), and $\Theta^{\mu}_{\nu}{}^{\mu}_{\lambda} := h^{\mu\alpha} \nabla^{\nu}_{\alpha} \Phi t_{\nu\lambda}$ representing the traceless gravitational field tensor associated with the Newtonian potential. Here $h^{\mu\nu}$ and $t_{\mu\nu}$, respectively, are the degenerate and mutually orthogonal spatial and temporal metrics with signatures (0 + + +) and (+000), representing the immutable chronogeometrical structure of the Newton-Cartan spacetime. They may be viewed as the ' $c \to \infty$ ' limits of the Lorentzian metric tensor field: $h^{\mu\nu} = \lim_{c \to \infty} (g^{\mu\nu}/c^2)$ and $t_{\mu\nu} = \lim_{c \to \infty} g_{\mu\nu}$. The conceptual superiority of this geometrization of Newtonian gravity is reflected in the trading of the two 'gauge-dependent' quantities Γ and Θ in favor of their gauge-independent sum Γ . Physically, it is the 'curved' connection Γ rather than any 'flat' connection Γ that can be determined by local experiments. Nither the potential Φ nor the 'flat' connection Γ has an independent existence; they exist

only relative to an arbitrary choice of a local inertial frame. It is worth noting that, unlike in both special and general theories of relativity, where the chronogeometrical structure of spacetime uniquely determines its inertio-gravitational structure, in Newton-Cartan theory these two structures are independently specified, subject only to the compatibility conditions $\nabla_{\alpha}h^{\beta\gamma}=0$ and $\nabla_{\alpha}t_{\beta\gamma}=0$. In fact, the connection Γ , as a solution of these compatibility conditions, is not unique unless a symmetry such as $R^{\alpha}_{\beta \cdot \delta}^{\gamma} = R^{\gamma}_{\delta \cdot \beta}^{\alpha}$ of the curvature tensor – capturing the 'curl-freeness' of the Newtonian gravitational filed – is assumed (here the indices are raised by the degenerate spatial metric $h^{\mu\nu}$). Further, although the two metric fields are immutable or non-dynamical in the sense that their Lie derivatives vanish identically,

$$\mathcal{L}_{\mathbf{x}}t_{\mu\nu} \equiv 0 \quad \text{and} \quad \mathcal{L}_{\mathbf{x}}h^{\mu\nu} \equiv 0 ,$$
 (5.3)

the connection field remains dynamical, $\mathcal{L}_x \Gamma_{\alpha \beta}^{\ \gamma} \neq 0$, since it is determined by the evolving distributions of matter. The generators $x = (t, x^a)$ of the 'isometry' group defined by the conditions (5.3), represented in an arbitrary reference frame, take the form (cf. Eq. (4.15))

$$t' = t + constant,$$

 $x'^{a} = O^{a}_{b}(t) x^{b} + c^{a}(t), \qquad (a,b = 1,2,3),$ (5.4)

where $O_h^a(t) \in SO(3)$ forms an orthonormal rotation matrix for each value of t (with Einstein's summation convention for like indices), and $\mathbf{c}(t) \in \mathbb{R}^3$ is an arbitrary time-dependent vector function. Physically, these transformations connect different observers in arbitrary (accelerating and rotating) relative motion.

With these physical motivations, the complete geometric set of gravitational field equations of the classical Newton-Cartan theory can be written as:

$$h^{\alpha\beta}t_{\beta\gamma} = 0, \quad \nabla_{\alpha}h^{\beta\gamma} = 0, \quad \nabla_{\alpha}t_{\beta\gamma} = 0, \quad \partial_{[\alpha}t_{\beta]\gamma} = 0,$$

$$R^{\alpha}_{\beta \cdot \delta} = R^{\gamma}_{\delta \cdot \beta},$$
and
$$R_{\mu\nu} + \Lambda t_{\mu\nu} = 4\pi G M_{\mu\nu},$$

$$(5.5a)$$

$$(5.5b)$$

$$R^{\alpha}_{\beta \cdot \delta}^{\gamma} = R^{\gamma}_{\delta \cdot \beta}^{\alpha}, \tag{5.5b}$$

and
$$R_{\mu\nu} + \Lambda t_{\mu\nu} = 4\pi G M_{\mu\nu}$$
, (5.5c)

where the first four equations specify the degenerate 'metric' structure and a set of torsion-free connections on the spacetime manifold \mathcal{M} , the fifth one picks out the Newton-Cartan connection from this set of generic possibilities, and the last one, with mass-momentum tensor $M_{\mu\nu} := \lim_{c\to\infty} T_{\mu\nu}$, relates spacetime geometry to matter in analogy with Einstein's field equations. Alternatively, one can recover this entire set of field equations (5.5) from Einstein's theory of gravity in the ' $c \to \infty$ ' limit (Künzle 1976, Ehlers 1981, 1986, 1991).

The only other field equation that is compatible with the structure (5.5) (Dixon 1975), but which cannot be recovered in the ' $c \to \infty$ ' limit of Einstein's theory, is

$$R^{\alpha\lambda}_{\ \ \gamma\delta} = 0 \tag{5.6}$$

(where, again, the index is raised by the degenerate spatial metric $h^{\lambda\sigma}$). It asserts the existence of absolute rotation in accordance with Newton's famous 'bucket experiment', and turns out to be of central importance for my argument against Penrose's experiment (cf. the next subsection). Without this extra field equation, however, there does not even exist a classical Lagrangian density for the Newton-Cartan system, let alone a Hamiltonian density or an unambiguous phase space. Despite many diligent attempts to construct a consistent Lagrangian density, the goal remains largely elusive, thanks to the intractable geometrical obstruction resulting from the degenerate 'metric' structure of the Newton-Cartan spacetime.

If, however, we take the condition (5.6) as an extraneously imposed but necessary field equation on the Newton-Cartan structure, then, after some tedious manipulations (cf. Christian 1997), we can obtain an unequivocal constraint-free phase space for the classical Newton-Cartan system coupled with Galileanrelativistic matter (Schrödinger fields). What is more, the restriction (5.6) also permits the existence of a family of local inertial frames in the Newton-Cartan structure (cf. the next subsection). Given such a local frame the inertial and gravitational parts of the Newton-Cartan connection-field can be unambiguously separated, as in the equation (5.2) above, and a non-rotating linear coordinate system may be introduced. Then, with some gauge choices appropriate for the earth-nucleus system of Penrose's experiment (recall that Penrose's experiment involves displacements of some 10¹⁵ nuclei), the relevant action functional (i.e., equation (4.3) of (Christian 1997)) takes the simplified form

$$\mathcal{I} = \int dt \int d\mathbf{x} \left[\frac{1}{8\pi G} \Phi \nabla^2 \Phi + \frac{\hbar^2}{2m} \delta^{ab} \partial_a \psi \partial_b \overline{\psi} + i \frac{\hbar}{2} \left(\psi \partial_t \overline{\psi} - \overline{\psi} \partial_t \psi \right) + m \overline{\psi} \psi \Phi \right], \tag{5.7}$$

where $\psi = \psi(\mathbf{x}_{CM}, \mathbf{x})$ is a complex Schrödinger field representing the composite earth-nucleus system, m is the reduced mass for the system, all spatial derivatives are with respect to the relative coordinate \mathbf{x} , and from now on the explicit reference to observer v on the top of the scalar Newtonian potential $\Phi(\mathbf{x})$ is omitted. Evidently, the convenient inertial frame I have chosen here is the CM-frame in which kinetic energy of the center-of-mass vanishes identically. In addition, one may also choose $\mathbf{x}_{CM} \equiv 0$ without loss of generality so that $\psi = \psi(\mathbf{x})$. Since the dynamics of the earth-nucleus system is entirely encapsulated in the function $\psi(\mathbf{x})$, it is sufficient to focus only on this \mathbf{x} -dependence of ψ and ignore the free motion of the center-of-mass. Needless to say that, since $m_{earth} \gg m_{nuleus}$, to an excellent approximation $m = m_{nuleus}$, and effectively the CM-frame is the laboratory-frame located at the center of the earth.

Extremization of the functional (5.7) with respect to variations of $\Phi(\mathbf{x})$ immediately yields the Newton-Poisson equation

$$\nabla^2 \Phi(\mathbf{x}) = -4\pi G m \,\overline{\psi}(\mathbf{x}) \,\psi(\mathbf{x}) \,, \tag{5.8}$$

which describes the manner in which a quantum mechanically treated particle bearing mass m gives rise to a 'quantized' gravitational potential $\Phi(\mathbf{x})$, thereby capturing the essence of Newtonian quantum gravity. On the other hand, extremization of the action with respect to variations of the matter field $\overline{\psi}(\mathbf{x})$ leads to the familiar Schrödinger equation for a quantum particle of mass m in the presence of an external field $\Phi(\mathbf{x})$:

$$i\hbar \frac{\partial}{\partial t} \psi(\mathbf{x}, t) = \left[-\frac{\hbar^2}{2m} \nabla^2 - m \Phi(\mathbf{x}) \right] \psi(\mathbf{x}, t) . \tag{5.9}$$

The last two equations may be reinterpreted as describing the evolution of a *single* particle of mass m interacting with its own Newtonian gravitational field. Then these coupled equations constitute a *nonlinear* system, which can be easily seen as such by first solving equation (5.8) for the potential $\Phi(\mathbf{x})$, giving

$$\Phi(\mathbf{x}) = G m \int d\mathbf{x}' \frac{\overline{\psi}(\mathbf{x}') \psi(\mathbf{x}')}{|\mathbf{x} - \mathbf{x}'|}, \qquad (5.10)$$

and then — by substituting this solution into equation (5.9) — obtaining the integro-differential equation (cf. equation 5.18 of (Christian 1997))

$$i\hbar \frac{\partial}{\partial t} \psi(\mathbf{x}, t) = -\frac{\hbar^2}{2m} \nabla^2 \psi(\mathbf{x}, t) - G m^2 \int d\mathbf{x}' \frac{\overline{\psi}(\mathbf{x}', t) \psi(\mathbf{x}', t)}{|\mathbf{x} - \mathbf{x}'|} \psi(\mathbf{x}, t).$$
 (5.11)

As alluded to at the end of subsection 4.3, Penrose has christened this equation 'Schrödinger-Newton equation', and regards the stationary solutions of it as the 'basic states' into which the quantum superpositions must reduce, within this Newtonian approximation of the full 'quantum gravity'.

As it stands, this equation is evidently a nonlinear equation describing a self-interacting quantum particle. However, if we promote ψ to a 'second-quantized' field operator $\widehat{\psi}$ satisfying (Christian 1997)

$$\left[\widehat{\psi}(\mathbf{x}), \widehat{\psi}^{\dagger}(\mathbf{x}')\right] = \widehat{\mathbb{1}} \delta(\mathbf{x} - \mathbf{x}')$$
(5.12)

at equal-times, then this equation corresponds to a linear system of many identical (bosonic) particles bearing mass m in the Heisenberg picture, with $\hat{\psi}$ acting as an annihilation operator in the corresponding Fock space. In particular, the properly normal-ordered Hamiltonian operator for the system now reads

$$\widehat{\mathbf{H}} = \widehat{\mathbf{H}}_{O} + \widehat{\mathbf{H}}_{I},$$
with
$$\widehat{\mathbf{H}}_{O} := \int d\mathbf{x} \ \widehat{\psi}^{\dagger}(\mathbf{x}) \left[-\frac{\hbar^{2}}{2m} \nabla^{2} \right] \widehat{\psi}(\mathbf{x})$$
and
$$\widehat{\mathbf{H}}_{I} := -\frac{1}{2} G m^{2} \int d\mathbf{x} \int d\mathbf{x}' \ \frac{\widehat{\psi}^{\dagger}(\mathbf{x}') \ \widehat{\psi}^{\dagger}(\mathbf{x}) \ \widehat{\psi}(\mathbf{x}) \ \widehat{\psi}(\mathbf{x}')}{|\mathbf{x} - \mathbf{x}'|},$$
(5.13)

which upon substitution into the Heisenberg equation of motion

$$i\hbar\frac{\partial}{\partial t}\widehat{\psi}(\mathbf{x},t) = \left[\widehat{\psi}(\mathbf{x},t),\widehat{\mathbf{H}}\right]$$
 (5.14)

yields an operator equation corresponding to (5.11). It is easy to show (Schweber 1961, 144) that the action of the Hamiltonian operator \widehat{H} on a multi-particle state $|\Psi\rangle$ is given by

$$\langle \mathbf{x}_1 \, \mathbf{x}_2 \, \dots \, \mathbf{x}_n | \, \widehat{\mathbf{H}} \, | \, \Psi \, \rangle = \left[-\frac{\hbar^2}{2m} \sum_{i=1}^n \nabla_i^2 \, - \, \frac{1}{2} \, G \, m^2 \sum_{\substack{i,j=1 \ i \neq j}}^n \frac{1}{|\mathbf{x}_i - \mathbf{x}_j|} \right] \langle \mathbf{x}_1 \, \mathbf{x}_2 \, \dots \, \mathbf{x}_n | \, \Psi \, \rangle \,, \tag{5.15}$$

which is indeed the correct action of the multi-particle Hamiltonian with gravitational pair-interactions. Put differently, since the Hamiltonian (5.13) annihilates any single-particle state, the particles no longer gravitationally self-interact. Thus, in a local inertial frame, the Newton-Cartan-Schrödinger system (Christian 1997) reduces, formally, to the very first quantum field theory constructed by Jordan and Klein (1927).

5.2. The inadequacy of Penrose's proposed experiment:

As noted above, the orthodox analysis carried out in the previous subsection is contingent upon the extraneously imposed field equation

$$R^{\alpha\lambda}_{\ \cdot\ \gamma\delta} = 0, \tag{5.16}$$

Eq. (5.6), without which the existence of even a classical Lagrangian density for the Newton-Cartan system seems impossible (cf. Christian 1997: subsection II C, subsection IV A, footnote 6). More significantly for our purposes, unless this extra condition prohibiting rotational holonomy is imposed on the curvature tensor, it is not possible to recover the Newton-Poisson equation (5.8),

$$\nabla^2 \Phi \left(\mathbf{x} \right) = -4\pi G \rho \left(\mathbf{x} \right) \,, \tag{5.17}$$

from the usual set of Newton-Cartan field equations (5.5) (which are obtained in the ' $c \to \infty$ ' limit of Einstein's theory) without any unphysical global assumption. Thus (5.16) embodies an essential discontinuity in the ' $c \to \infty$ ' limit between the gravitational theories of Einstein and Newton (cf. Figure 1), and without it the Schrödinger-Newton equation (5.9) is not meaningful.

Let us look at this state of affairs more closely (cf. Misner *et al.* 1973, 294-295, Ehlers 1981, 1986, 1991, 1997). The only nonzero components of the connection-field corresponding to the set of field equations (5.5) (and the coordinate transformations (5.4)) are

$$\Gamma_{0\ 0}^{\ a} =: -g^{a} \quad \text{and} \quad \Gamma_{0\ a}^{\ b} = O_{\ c}^{\ b} \dot{O}_{\ a}^{c} := h^{bc} \varepsilon_{acd} \Omega^{d}.$$
 (5.18)

With respect to a coordinate system, the spatial vector fields $\mathbf{g}(\mathbf{x}, t)$ and $\mathbf{\Omega}(\mathbf{x}, t)$ play the part of gravitational acceleration and Coriolis angular velocity, respectively, and the field equations (5.5) reduce to the set

$$\nabla \cdot \mathbf{\Omega} = 0, \qquad \nabla \times \mathbf{g} + 2\dot{\mathbf{\Omega}} = 0,$$

$$\nabla \times \mathbf{\Omega} = 0, \qquad \nabla \cdot \mathbf{g} - 2\mathbf{\Omega}^2 = 4\pi G \rho,$$
(5.19)

where \mathbf{g} and $\mathbf{\Omega}$ in general depend on both \mathbf{x} and t (and I have set $\Lambda=0$ for simplicity). It is clear from this set that the recovery of the Newton-Poisson equation – and hence the reduction to the *strictly-Newtonian* theory – is possible if and only if a coordinate system exists with respect to which $\mathbf{\Omega}=0$ holds. This can be achieved if $\mathbf{\Omega}$ is spatially constant – i.e., depends on time only. And this is precisely what is ensured by the extra field equation (5.16), which asserts that the parallel-transport of spacelike vectors is path-independent. Given this condition, the coordinate system can be further specialized to a nonrotating one, with $\Gamma_0{}^b{}_a=0$, and the connection coefficients can be decomposed as in equation (5.2), with $\mathbf{g}:=-\nabla\Phi$.

This entire procedure, of course, may be sidestepped if we admit only asymptotically flat spacetimes. With such a global boundary condition, the restriction (5.16) on the curvature tensor becomes redundant (Künzle 1972, Dixon 1975). However, physical evidence clearly suggests that we are *not* living in an 'island universe' (cf. Penrose 1996, 593-594) – i.e., universe is not 'an island of matter surrounded by emptiness' (Misner *et al.* 1973, 295). Therefore, a better procedure of recovering the Newtonian theory from Einstein's theory is not to impose such a strong and unphysical global boundary condition, but, instead, to require that only the weaker condition on the curvature tensor, (5.16), is satisfied. For, this weaker condition is quite sufficient to recover the usual version of Newton's theory with gravitation as a force field on a flat, non-dynamical, a priori spacetime structure, and guarantees existence of a class of inertial coordinate systems not rotating with respect to each other; i.e., the condition suppresses time-dependence of the rotation matrix $O_b^a(t)$ (as a result of the restriction $\Gamma_0^b{}_a = 0$), and reduces the transformation law (5.4) to

$$t \longrightarrow t' = t + constant,$$

$$x^a \longrightarrow x'^a = O^a_b x^b + c^a(t), \qquad (a,b = 1,2,3). \qquad (5.20)$$

Note that, unlike the asymptotic-flatness imposing condition $\lim_{|\mathbf{x}|\to\infty} \Phi(\mathbf{x}) = 0$, the weaker condition (5.16) does not suppress the arbitrary time-dependence of the function $c^a(t)$ – i.e., (5.16) does not reduce $c^a(t)$ to $v^a \times t$ as in the Galilean transformation (4.15) above. Consequently, the gravitational potential Φ in the resultant Newtonian theory remains nonunique (Misner *et al.* 1973, 295), and, under the diffeomorphism corresponding to the transformation (5.20), transforms (actively) as

$$\Phi(\mathbf{x}) \longrightarrow \Phi'(\mathbf{x}) = \Phi(\mathbf{x}) - \ddot{\mathbf{c}} \cdot \mathbf{x}. \tag{5.21}$$

Let us now go back to Penrose's hypothesis on the mechanism underlying quantum state reduction discussed in the subsection 4.2 above, and retrace the steps of that subsection within the present strictly-Newtonian scenario. As before, although here $h^{\mu\nu}$ and $t_{\mu\nu}$ would serve as 'individuating fields' (cf. section 2) allowing pointwise identification between two different spacetimes, due to the transformation law (5.20) there would appear to be an ambiguity in the notion of time-translation operator analogous to Eq. (4.16),

$$\frac{\partial}{\partial \mathbf{x}'^{a}} = \frac{\partial}{\partial \mathbf{x}^{a}} \quad \text{but} \quad \frac{\partial}{\partial t} \longrightarrow \frac{\partial}{\partial t'} = \frac{\partial}{\partial t} - \dot{\mathbf{c}}^{a}(t) \frac{\partial}{\partial \mathbf{x}^{a}}, \tag{5.22}$$

when superpositions involving two such different spacetimes are considered. However, I submit that this 'ambiguity' in the present – essentially Newtonian – case is entirely innocuous. For, in the strictly Newtonian theory being discussed here, where a 'spacetime' now is simply a flat structure 'plus' a gravitational potential $\Phi(\mathbf{x})$ as in equation (5.2), one must consider (5.22) together with the transformation (5.21). But the Schrödinger-Newton equation (5.9) – which is the appropriate equation here – happens to be covariant under such a concurrent transformation⁶, and retains the original form

$$i\hbar \frac{\partial}{\partial t'} \psi'(\mathbf{x}, t) = \left[-\frac{\hbar^2}{2m} \nabla'^2 - m \Phi'(\mathbf{x}) \right] \psi'(\mathbf{x}, t)$$
 (5.23)

(Rosen 1972, cf. also Christian 1997) with the following (active) transformation of its solution (if it exists):

$$\psi(\mathbf{x}, t) \longrightarrow \psi'(\mathbf{x}, t) = e^{if(\mathbf{x}, t)} \psi(\mathbf{x}, t). \tag{5.24}$$

What is more (cf. Kuchař 1980, 1991), due to inverse relation between transformations on the function space and transformations (5.20) on coordinates, equation (5.24) implies

$$\psi'(\mathbf{x}',t') = \psi(\mathbf{x},t) . \tag{5.25}$$

Better still: under simultaneous gauge transformations (5.21), (5.22) and (5.24), the Lagrangian density of the action (5.7) remains invariant except for a change in the spatial boundary term, which of course does not contribute to the Euler-Lagrange equations (5.8) and (5.9). Thus, the entire Schrödinger-Newton theory is unaffected by these transformations, implying that it is independent of a particular choice of reference frame represented by $\frac{\partial}{\partial t}$ out of the whole family given in (5.22). It should be noted, however, that here, as in any such demonstration of covariance, all variations $\delta\Phi$ of the Newtonian potential is assumed to vanish identically at the spatial boundary.

That is to say, the new solution of the Schrödinger-Newton equation expressed in the new coordinate system is *exactly* equal to the old solution expressed in the old coordinate system – the new value of the ψ -field, as measured at the transformed spacetime point, is numerically the same as its old value measured at the original spacetime point. Now consider a superposition involving two entirely different strictly-Newtonian 'spacetimes' in the coordinate representation analogous to the 'superposition' (4.19) discussed in section 4,

$$\langle \mathbf{x} | \mathcal{X}(t) \rangle = \lambda_2 \Psi_2(\mathbf{x}, t) + \lambda_3 \Psi'_3(\mathbf{x}', t') ,$$
 (5.26)

where unprimed coordinates correspond to one spacetime and the primed coordinates to another⁷. *Prima facie*, in accordance with the reasonings of section 4, such a superposition should be as unstable as Eq. (4.19). However, in the present strictly-Newtonian case, thanks to the relation (5.25), the physical state represented by (5.26) is equivalent to the superposed state

$$\langle \mathbf{x} | \mathcal{X}(t) \rangle = \lambda_2 \Psi_2(\mathbf{x}, t) + \lambda_3 \Psi_3(\mathbf{x}, t) . \tag{5.27}$$

And there is, of course, nothing unstable about such a superposition in this strictly-Newtonian domain. Consequently, for such a superposition, $\Delta E \equiv 0$, and hence its life-time $\tau \sim \infty$ (cf. Eq. (4.21)).

Thus, as long as restriction (5.16) on the curvature tensor is satisfied – i.e., as long as it is possible to choose a coordinate system with respect to which $\Gamma_{0a}^{\ b} = 0$ holds for each spacetime, the Penrosetype instability in quantum superpositions is non-existent (a conclusion not inconsistent with the results of (Christian 1997)). Put differently, given $\Gamma_{0a}^{b} = 0$, the Penrose-type obstruction to stability of superpositions is sufficiently mitigated to sustain stable quantum superpositions. In physical terms, since (5.16) postulates the existence of 'absolute rotation', the superposition (5.26) is perfectly Penrose-stable as long as there is no relative rotation involved between its two components. On the other hand, if there is a relative rotation between the two components of (5.26) so that $\Gamma_{0a}^{b} = 0$ does not hold for both spacetimes, than it is not possible to analyze the physical system in terms of the strictly-Newtonian limit of Einstein's theory, and, as a result, the 'superposition' (5.26) would be Penrose-unstable. Unfortunately, neither in Penrose's original experiment (1998), nor in the version discussed in subsection 4.5 above, is there any relative rotation between two components of the superposed mass distributions. In other words, in both cases $\Gamma_{0a}^{b} = 0$ holds everywhere, and hence no Penrose-type instability should be expected in the outcome of these experiments. (Incidentally, among the known solutions of Einstein's field equations, the only known solution which has a genuinely Newton-Cartan limit – i.e., in which Ω is not spatially constant, entailing that it cannot be reduced to the strictly-Newtonian case with $\Gamma_{0a}^{b} = 0$ – is the NUT spacetime (Ehlers 1997)).

5.3. More adequate experiments involving relative rotations:

It is clear from the discussion above that, in order to detect Penrose-type instability in superpositions, what we must look for is a physical system for which the components $\Gamma_0{}^b{}_a$ of the connection field, in addition to the components $\Gamma_0{}^a{}_0$, are meaningfully non-zero. Most conveniently, there exists extensive theoretical and experimental work on just the kind of physical systems we require.

The first among these systems involves 'macroscopic' superpositions of two screening currents in r.f.-SQUID rings, first proposed by Leggett almost two decades ago (Leggett 1980, 1984, 1998, Leggett and Garg 1985). An r.f.-SQUID ring consists of a loop of superconducting material interrupted by a thin Josephson tunnel junction. A persistent screening current may be generated around the loop in response to an externally applied magnetic flux, which obeys an equation of motion similar to that of a particle moving in a one dimensional double-well potential. The thus generated current in the ring would be equal in magnitude in both wells, but opposite in direction. If dissipation in the junction and decoherence due to environment are negligible, then the orthodox quantum analysis predicts coherent oscillations between the two distinct flux states, and, as a result, a coherent superposition between a large number of electrons flowing around the ring

⁷ Of course, since the Schrödinger-Newton equation is a non-linear equation, its more adequate (orthodox) quantum-mechanical treatment is the one given by equations (5.12)–(5.15) of the previous subsection. My purpose here, however, is simply to parallel Penrose's argument of instability in quantum superpositions near the Planck mass.

in opposite directions – clockwise or counterclockwise – is expected to exist, generating a physical situation analogous to the one in Eq. (4.19) or (5.26) above. Most importantly for our purposes, since there would be relative rotation involved between the currents in the two possible states, owing to the Lense-Thirring fields (Lense and Thirring 1918, Ciufolini *et al.* 1998) of these currents, the connection components $\Gamma_0{}^b{}_a$, in addition to the components $\Gamma_0{}^a{}_0$, will be nonzero. And this will unambiguously give rise to a Penrose-type instability at an appropriate mass scale – say roughly around 10^{21} electrons. The number of electrons in the SQUID ring in an actual experiment currently under scrutiny in Italy (Castellano *et al.* 1996) is only of the order of 10^{15} , but there is no reason for a *theoretical* upperbound on this number.

It should be noted that Penrose himself has briefly considered the possibility of a Leggett-type experiment to test his proposal (1994b, 343). Recently, Anandan (1998) has generalized Penrose's expression for ΔE to arbitrary connection fields (cf. footnote 5), which allows him to consider connection components other than $\Gamma_0{}^a{}_0$, in particular the components $\Gamma_0{}^b{}_a$, and suggest a quantitative test of Penrose's ansatz via Leggett's experiment. What is novel in my own endorsement of this suggestion is the realization that Leggett-type experiments belong to a class of experiments – namely, the class involving $\Gamma_0{}^b{}_a \neq 0$ – which is the only class available within the nonrelativistic domain to unequivocally test Penrose's proposal.

A second more exotic physical system belonging to this class of experiments is a superposition of two vortex states of an ultracold Bose-Einstein Condensate (BEC), currently being studied by Cirac's group in Austria among others (Cirac et al. 1998, Dum et al. 1998, Butts and Rokhsar 1999). Again, owing to the Lense-Thirring fields of such a slowly whirling BEC (clockwise or counterclockwise), a Penrose-type instability can in principle be detected at an appropriate mass scale.

Finally, let me point out that the analysis of this section has opened up an exciting new possibility of empirically distinguishing Penrose's scheme from other (ad hoc) theories of gravity-induced state reduction (e.g., Ghirardi et al. 1990), with the locus of differentiation being the connection components $\Gamma_0{}^b{}_a$. There is nothing intrinsic in such ad hoc theories that could stop a state from reducing when these connection components are zero – e.g., for the experiment described in subsection 4.5 above these theories predict reduction at an appropriate scale, whereas Penrose's scheme, for the reasons explicated above, does not.

6. Concluding Remarks:

Notwithstanding the importance of partial reservations levelled against Penrose's proposed experiment in the previous section, it should be clear that my criticism has significance only in the strictly-Newtonian domain. The classical world, of course, is not governed by Galilean-relativistic geometries, but by general-relativistic geometries. Accordingly, the true domain of the discussion under consideration must be the domain of full 'quantum gravity'. And, reflecting on this domain, I completely share Penrose's sentiments that "our present picture of physical reality, particularly in relation to the nature of time, is due for a grand shake up" (1989, 371) (similar sentiments, arrived at from quite a different direction, are also expressed by Shimony (1998)). The incompatibility between the fundamental principles of our two most basic theories – general relativity and quantum mechanics – is so severe that the unflinching orthodox view maintaining a status quo for quantum superpositions – including at such a special scale as the Planck scale – is truly baffling. As brought out in several of the essays in this collection and elaborated by myself in section 4 above, the conflict between the two foundational theories has primarily to do with the axiomatically presupposed fixed causal structure underlying quantum dynamics, and the meaninglessness of such a fixed, non-dynamical, background causal structure in the general relativistic picture of the world. The orthodox response to the conflict is to hold the fundamental principles of quantum mechanics absolutely sacrosanct at the price of severe compromises with those of Einstein's theory of gravity. For example, Banks, one of the pioneers of the currently popular M-theory program, has proclaimed (1998b): "... it seems quite clear that the fundamental rules of [Mtheory] will seem outlandish to anyone with a background in ... general relativity. ... At the moment it appears that the only things which may remain unscathed are the fundamental principles of quantum mechanics." In contrast, representing a view of growing minority, Penrose has argued for a physically more meaningful evenhanded approach in which even the superposition principle is not held beyond reproach at all scales. It certainly requires an extraordinary leap of faith in quantum mechanics (a leap, to be precise, of some seventeen orders of magnitude in the length scale!) to maintain that the Gordian knot – the conflict between our two most basic theories – can be cut without compromising the superposition principle in some manner. My own feeling, heightened by Penrose's tenacious line of reasoning, is that such a faith in quantum mechanics could turn out to be fundamentally misplaced, as so tellingly made plain by Leggett (1998):

"Imagine going back to the year 1895 and telling one's colleagues that classical mechanics would break down when the product of energy and time reached a value of order 10^{-34} joule seconds. They would no doubt respond gently but firmly that any such idea must be complete nonsense, since it is totally obvious that the structure of classical mechanics cannot tolerate any such characteristic scale!"

Indeed, one often comes across similar sentiments with regard to the beautiful internal coherence of quantum formalism. However, considering the extraordinary specialness of the Planck scale, I sincerely hope that our 'quantum' colleagues are far less complacent than their 'classical' counterparts while harbouring the 'dreams of a final theory'.

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